

Core-collapse supernova simulations and the formation of neutron stars, hybrid stars, and black holes

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ABSTRACT

We investigate observable signatures of a first-order quantum chromodynamics (QCD) phase transition in the context of core collapse supernovae. To this end, we conduct axially symmetric numerical relativity simulations with multi-energy neutrino transport, using a hadron-quark hybrid equation of state (EOS). We consider four non-rotating progenitor models, whose masses range from 9.6 to 70 M_{\odot} . We find that the two less massive progenitor stars (9.6 and 11.2 M_{\odot}) show a successful explosion, which is driven by the neutrino heating. They do not undergo the QCD phase transition and leave behind a neutron star (NS). As for the more massive progenitor stars (50 and 70 M_{\odot}), the proto-neutron star (PNS) core enters the phase transition region and experiences the second collapse. Because of a sudden stiffening of the EOS entering to the pure quark matter regime, a strong shock wave is formed and blows off the PNS envelope in the 50 M_{\odot} model. Consequently the remnant becomes a quark core surrounded by hadronic matters, leading to the formation of the hybrid star. However for the 70 M_{\odot} model, the shock wave cannot overcome the continuous mass accretion and it readily becomes a black hole. We find that the neutrino and gravitational wave (GW) signals from supernova explosions driven by the hadron-quark phase transition are detectable for the present generation of neutrino and GW detectors. Furthermore, the analysis of the GW detector response reveals unique kHz signatures, which will allow us to distinguish this class of supernova explosions from failed and neutrino-driven explosions.

Keywords: Supernova dynamics (1664), Compact objects (288), High energy astrophysics (739), Supernova neutrinos (1666), Gravitational waves (678), Hydrodynamics (1963)

1. INTRODUCTION

The final fate of massive stars is the moment of birth of diverse compact stars. Massive stars heavier than $\sim 8 M_{\odot}$ undergo gravitational collapse of their iron core once their core mass reaches the Chandrasekhar limit. Thereafter part of or whole the star collapses and becomes a compact remnant star: a neutron star (NS) or a black hole (BH) or an *exotic* star. The remnant type is determined mostly from the progenitor mass (Heger et al. 2003). Generally speaking, more massive stars tend to leave behind BHs and conversely less massive stars have more chances of leaving NSs (Fryer et al. 1999).

The gravitational collapse releases huge amount of gravitational binding energy, part of which is deposited into the stellar envelop. Depending on the progenitor star, the energy deposition can trigger the stellar explosion called core-collapse

supernova (CCSN). The explosion dynamics regulates how much mass is blown off, or equivalently how much mass accretes on the remnant star. Therefore, the detailed explosion dynamics including the success or failure of the explosion are the keys to understand the diversity of remnant compact stars.

Two major explosion mechanisms are currently proposed. The first one relies on the complex interplay between neutrinos and stellar mantle, called neutrino-driven explosion (NDE). It is considered to account for the vast majority of observed CCSNe, showing an explosion energy of $\sim 10^{51}$ erg (for recent reviews, see Janka et al. 2016; Radice et al. 2018; Müller 2020). The other is called magnetorotational explosion (MRE) that takes place if the progenitor star rotates sufficiently fast and is also magnetized (see, Mösta et al. 2018; Obergaulinger & Aloy 2021; Kuroda 2021; Bugli et al. 2021, for recent three-dimensional studies). Due to its efficient energy conversion mechanism, the MRE mechanism may explain some exceptional explosive events called hypernovae (Iwamoto et al. 1998; Nomoto et al. 2006), which exhibit one order of magnitude larger explosion energy than the canonical

events. Roughly speaking, the NDE and MRE are currently presumed to explain CCSNe of less and more massive stars, respectively.

Recently another explosion mechanism gains more attention in the context of both the explosion dynamics and formation process of neutron stars composed of exotic matter at their interiors (c.f. Fischer et al. 2018; Zha et al. 2020). When the energy density becomes sufficiently high then nuclear, in general hadronic, matter undergoes a phase transition to the deconfined quark-gluon plasma. Therefore, the solution of quantum chromodynamic (QCD)—the theory for strong interactions with quarks and gluons as the degrees of freedom—predicts a smooth cross-over transition at a pseudocritical temperature of in the range of 150–160 MeV Bazavov et al. (2014); Borsányi et al. (2014); Bazavov et al. (2019), however, only at vanishing baryon density and hence inherent to astrophysical applications of compact stellar objects such as neutron stars and supernovae. At high density phenomenological quark matter models have long been employed in astrophysical studies, based on the common two-phase approach with separate hadronic and quark matter phases. Such phase transition construction results not only in the transition to the third family of compact stars, known as *hybrid* stars, as an extension to the neutron stars second family, but also in the yet incompletely understood question about the nature of the QCD phase transition at high baryon density (c.f. Annala et al. 2020, and references therein). It is likely that low- and intermediate-mass NSs, having masses of $\sim 1.4 M_{\odot}$, can be explained solely by hadronic equations of state (EOS). More massive NSs with $\sim 2 M_{\odot}$, whose presence is observationally confirmed at high accuracy (c.f. Demorest et al. 2010; Antoniadis et al. 2013; Cromartie et al. 2020), feature the highest baryon densities encountered in the universe. It is a valid question to ask if a QCD transition occurs at the densities encountered at the interior of such massive pulsars (c.f. Kurkela et al. 2014, and references therein). Furthermore, if confirmed, this is likely to affect the understanding of the core-collapse supernova phenomenology.

This is the origin of the third explosion mechanism, which utilizes the difference in the gravitational binding energy between NSs and hybrid stars. As a consequence, once the central energy density of the PNS enters the phase transition region, the PNS core migrates to a new state, liberating considerable amount of gravitational potential energy. The latter is directly related to the latent heat, as a consequence of the phase transition construction within the two-phase approach. (for a recent review, see Baym et al. 2018). Assuming a sufficiently strong first order transition, previous CCSN simulations showed a possible explosion scenario (Sagert et al. 2009; Nakazato et al. 2013; Fischer et al. 2018; Zha et al. 2020), which relies on the energy liberation at the moment of the phase transition. Under such conditions, the

PNS core favors a dynamical collapse, rather than a smooth transition, into the hybrid star branch. When the infalling matter collides with a solid surface of the quark core, the dynamical collapse can be halted depending on the PNS mass. Then it leads to a formation of a strong shock wave, which may blow off the outer envelope. This is the mechanism of the explosion driven by the QCD phase transition (hereafter, phase transition driven explosion, PDE).

It is intriguing to explore the progenitor mass dependence on the PDE dynamics and also on the final remnant property. One can naively expect that progenitors on the less massive side (say $\lesssim 20 M_{\odot}$) may explode as CCSNe (e.g., Müller et al. 2013; Takiwaki et al. 2016; Vartanyan et al. 2019), probably by the standard NDE mechanism, and leave NSs. This is because, for such low mass stars, the PDE may not be relevant, as the PNS interior unlikely becomes so dense enough to reach the QCD phase transition before the NDE. While the more massive stars ($\gtrsim 70 M_{\odot}$) tend to become failed CCSNe and leave BHs (Kuroda et al. 2018; Shibagaki et al. 2021). Although the progenitor mass threshold, above which they are expected to end up with BH formations, is yet to be understood, some intermediate-mass stars, e.g., $\sim 40 - 50 M_{\odot}$ at their zero-age main sequence (ZAMS), are reported to form BHs in the standard neutrino heating mechanism (irrespective of the spatial dimension: Liebendörfer et al. 2004; Sumiyoshi et al. 2007; Chan et al. 2018; Pan et al. 2018). However, with the aid of the QCD phase transition, these intermediate-mass stars could turn into successful CCSNe with the remnant being hybrid stars (Fischer et al. 2018). Progenitor mass dependence of phase transition is studied by Fischer et al. (2020a); Zha et al. (2021), though only in spherically symmetric models.

It is noteworthy that, contrary to the NDE case, the PDE mechanism works even in the spherically symmetric model. There is a consensus that the successful CCSN explosion driven by neutrino heating cannot be achieved under the spherical symmetry (1D) (Liebendörfer et al. 2001; Sumiyoshi et al. 2005), except for a few less massive star models (with $\lesssim 10 M_{\odot}$, Kitaura et al. 2006). This is because the inefficient neutrino heating in the absence of lateral hydrodynamic motions in 1D. The mechanism of PDE, on the contrary, does not rely on the neutrino heating, but mostly on the strongness of hydrodynamic bounce of quark core, i.e. on the QCD phase transition. Therefore most of previous PDE studies are reported using 1D models (Sagert et al. 2009; Nakazato et al. 2013; Fischer et al. 2018, 2020b) and the only multi dimensional (multi-D) model reported so far is by Zha et al. (2020). However, their study focuses mainly on the impacts of phase transition on the gravitational wave (GW) emission using a particular progenitor star with $12 M_{\odot}$, and the hydrodynamic evolution of multi-D PDE model is thus currently not well studied.

There are additional remarkable aspects of the PDE as follows: its relatively large explosion energy of a few times 10^{51} erg (Fischer et al. 2018) and a possible mass window of its progenitor stars. For the latter, the intermediate-mass stars with $\sim 40\text{-}50 M_{\odot}$ at ZAMS stage could be the potential candidates of the PDE as less massive stars ($\lesssim 20 M_{\odot}$) may explode solely by the neutrino heating mechanism, while too massive stars can directly collapse into a BH (Heger et al. 2003). According to the stellar evolution models, such intermediate-mass stars experience severe mass loss during their evolutionary stages (e.g., Umeda & Nomoto 2008) and would produce a dense circumstellar medium (CSM). From these facts, the PDE might be able to account for a subclass of CCSNe: the CSM-interacting superluminous SN (hereafter SLSN). This is because the observed light curves of SLSNe, such as SN 2006gy (Smith et al. 2007; Moriya et al. 2013) or SN 2008fz (Drake et al. 2010), can be well modeled by slightly larger explosion energy than normal type IIp SNe (Hamuy 2003), a few times 10^{51} erg, and a dense CSM, with which the SN explosion interacts. We should, however, note that the origin of such superluminous transient events are still under debate. While the above-mentioned interaction between type II SN ejecta and a dense CSM is indeed one plausible mechanism, other possible scenarios are also proposed to explain the light curves of SLSNe: such as by pulsational pair instability supernovae (Woosley et al. 2007) or type Ia SNe in CSM (Jerkstrand et al. 2020).

In this study, we perform 2D axially symmetric numerical simulations of core collapse of massive stars and explore the possible impacts of the phase transition on the PNS evolution and the remnant property. We particularly focus on the PDE dynamics in multi-D model. To this end, we select four massive progenitor stars ranging from $9.6 M_{\odot}$ to $70 M_{\odot}$ and use one particular EOS, which considers the first order phase transition from hadronic matter to quark-gluon plasma (for details, see Fischer et al. 2018; Bastian 2021). The employed supernova model is based on numerical relativity and solves the general relativistic neutrino radiation hydrodynamics equations together with the two-moment (M1) neutrino transport equations (details can be found in Kuroda et al. 2016b). We use up-to-date neutrino opacities following Kotake et al. (2018). Our results clearly show a diversity of postbounce evolution for various progenitor masses. We also discuss their multi-messenger signals: gravitational waves (GWs) and neutrino signals.

This paper is organized as follows. Section 2 starts with a concise summary of our GR radiation-hydrodynamic (RHD) scheme and also describe the initial setup of the simulation. The main results and detailed analysis of the effects of phase transition are presented in Section 3. In Section 4, we discuss multi-messenger signals obtained from computed models. We summarize our results and conclude in Section 5.

Note that cgs units are used throughout the paper. Greek indices run from 0 to 3 and Latin indices from 1 to 3, except ν and ε that denote neutrino species and energy, respectively.

2. METHOD AND INITIAL MODELS

In our full GR RHD simulations, we solve the evolution equations of geometrical variables and the relativistic hydrodynamics together with the energy-dependent neutrino radiation transport. The basic part in solving the evolution equations is the same as our previous full 3D studies in the Cartesian coordinates (Kuroda et al. 2016b, 2020). The only difference is that we employ the so-called Cartoon method (Shibata 2000; Alcubierre et al. 2001), which allows us to evolve the axisymmetric geometrical variables in the 3D Cartesian coordinates. On the other hand, we use the cylindrical coordinates for the RHD part, which is more suitable for satisfying the conservation laws, and solve the conservation equations. The physical domain is on $y = 0$ plane and the cell centers of the Cartoon meshes and cylindrical meshes overlap each other. In addition, all geometrical, hydrodynamical, and radiation variables are defined at cell centers.

We follow the collapse and post-bounce evolution of non-rotating progenitor stars with $9.6 M_{\odot}$, $11.2 M_{\odot}$, $50 M_{\odot}$, and $70 M_{\odot}$. The progenitor with $9.6 M_{\odot}$ (hereafter $z9.6$) is a star with zero-metallicity and is considered to be in the lowest mass range of progenitor stars of CCSNe. Its structural feature is a steep density gradient above the iron core that enables a successful explosion even in 1D models (Radice et al. 2017). Another low mass progenitor model with $11.2 M_{\odot}$ ($s11.2$) is a widely used progenitor model with solar metallicity and taken from Woosley et al. (2002). The progenitor model with $50 M_{\odot}$ ($s50$) of Umeda & Nomoto (2008) is a blue-supergiant star with solar-metallicity and was used in the previous study of Fischer et al. (2018) to explore a possible scenario of superluminous SN triggered by the QCD phase transition. The heaviest progenitor star among our set of progenitors is $70 M_{\odot}$ ($z70$) with zero metallicity (Takahashi et al. 2014). This progenitor star may form a BH, as was explored by Kuroda et al. (2018) and Shibagaki et al. (2021) and is a suitable model to investigate if the QCD phase transition can turn the failed SN model into a successful one.

In the present study we implement the EOS of Fischer et al. (2018). It features a phase transition from the DD2 nuclear relativistic mean field model of Typel et al. (2010), with density-dependent meson-nucleon couplings, to the string-flip quark matter EOS. The latter is based on a novel relativistic density functional (RDF) approach, which is build on a mechanism for (de)confinement through a string-like potential which is divergent at vanishing density. Furthermore, linear and higher-order repulsive vector interactions are included in the quark matter phase following Benić et al. (2015), which give rise to additional pressure contributions at high

density. The RDF model has been developed originally for cold, neutron star studies by Kaltenborn et al. (2017), which was extended by Fischer et al. (2018) and Bastian (2021) to finite temperatures and arbitrary isospin asymmetry, in order to be applicable to astrophysical studies of core-collapse supernovae in Fischer et al. (2018) and Fischer et al. (2020b) as well as binary neutron star mergers in Bauswein et al. (2019). Finite temperature and isospin asymmetry aspects are particularly important since the conditions for the onset of the phase transition have a strong temperature and charge fraction dependence (c.f. Fig. 1 and the phase diagram in Fig. 7 of Fischer 2021). Here, we employ the string-flip parametrization denoted as DD2F-RDF 1.2 (see Table I in Bastian 2021). It features densities for the onset of the phase transition of $3.25 \times \rho_{\text{sat}}^1$ at zero temperature ($T = 0$) and $1.15 \times \rho_{\text{sat}}$ at $T \simeq 50$ MeV (see Table 1 and Fig. 1 in Fischer 2021). The maximum mass of DD2F-RDF 1.2 for cold and non-rotating hybrid stars of $2.16 M_{\odot}$ is in agreement with the most precisely known massive pulsar of $2.08 \pm 0.07 M_{\odot}$ for the PSR J0740+6620 pulsar obtained by the analysis of Antoniadis et al. (2013), which has been recently revisited by Fonseca et al. (2021). We note further that the DD2F nuclear relativistic mean field EOS has been modified from its original DD2 version in order to be consistent with the supersaturation density constraint deduced from the heavy-ion collision data analysis of the elliptic flow by Danielewicz et al. (2002). For the phase transition construction from DD2F to RDF, a Maxwell-like construction is applied (further details can be found in Bastian 2021).

The 2D axially symmetric computational domain extends to 1.5×10^4 km, for $z70$ and $s50$, and to 7.5×10^3 km, for $s11.2$ and $z9.6$, from the center. Reason of the smaller computational domain for $s11.2$ and $z9.6$ is simply because that our EOS tabulation does not cover the required low density region. In the computational domain, nested boxes with 11(10) refinement levels are embedded for $z70$ and $s50$ ($s11.2$ and $z9.6$), and each nested box contains 64×64 cells so that the finest resolution at the center becomes ~ 115 m (a detailed resolution study can be found in Appendix A). The neutrino energy space ε logarithmically covers from 1 to 300 MeV with 20 energy bins.

In this study, we use the up-to-date neutrino rates of Kotake et al. (2018). The electron capture rate on heavy nuclei is the most elaborate one following Juodagalvis et al. (2010). Furthermore, we also take into account the following corrections: inelastic contributions and weak magnetism corrections of Horowitz (2002) for the charged-current neutrino emission and absorption processes involving the unbound

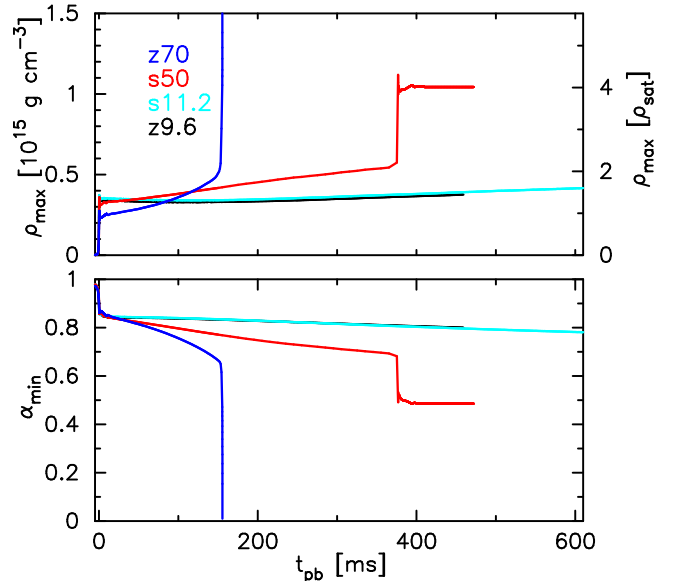


Figure 1. *Top panel:* evolution of the maximum rest mass density for all models in units of $10^{15} \text{ g cm}^{-3}$ (left axis) and of ρ_{sat} (right axis). *Bottom panel:* evolution of the central lapse function.

nucleons as well as for the neutral-current neutrino nucleon scattering processes; a correction for the effective nucleon mass (Reddy et al. 1999); the quenching of the axial-vector coupling constant (Carter & Prakash 2002; Fischer 2016); many-body and virial corrections of Horowitz et al. (2017) for the neutral-current neutrino nucleon scattering rates; effective strangeness contributions reducing the axial-vector coupling constant following Horowitz (2002), also here for both charged-current neutrino emission and absorption processes involving the unbound nucleons as well as for the neutral-current neutrino nucleon scattering processes. Following Melson et al. (2015a), we use $g_A^s = -0.1$ (Hobbs et al. 2016) for the effective strangeness contributions in this work, where this value increases(decreases) the cross section for νp (νn)-scattering.

3. RESULTS

In this section, we present results of our simulations. We begin with a description of overall dynamics of all models. After that, we focus on a particular model $s50$ and explain its second collapse and bounce in detail.

3.1. Collapse of stellar and PNS core

Figure 1 shows the evolution of maximum density and minimum lapse function for all models, from which we can gain a first glance of overall dynamics. After the central density exceeds nuclear saturation density, the stellar core collapse halts and the core experiences the first core bounce due to stiffening of the nuclear EOS. We define the moment of core bounce, $t_{\text{pb}} = 0$ when the maximum rest mass density is reached. It corresponds to the following values for the maximum, central

¹The saturation density has a value of $\rho_{\text{sat}} = 2.5 \times 10^{14} \text{ g cm}^{-3}$ or equivalent in nuclear units 0.15 fm^{-3} .

density of $\rho_{\max} = 3.81, 3.92, 3.62$ and $2.73 \times 10^{14} \text{ g cm}^{-3}$ for models $z9.6, s11.2, s50$ and $z70$, respectively. Roughly speaking, therefore, the maximum density at bounce has an inverse relation with the progenitor mass. This is related to the temperatures obtained in these simulations, which are highest(lowest) for the more(less) massive progenitors, i.e. the thermal pressure additionally supports the halt of the stellar core collapse and the subsequent core bounce.

After the first bounce, the maximum density again increases with time in all models due to the continuous mass accretion prior to the possible onset of a shock revival. While the less massive progenitor stars, $z9.6$ and $s11.2$, show only a slight increase compared to the other two more massive stars, the model $z70$ shows the most rapid increase rate, which features the highest mass accretion rate. Consequently, $z70$ reaches the highest central temperatures and luminosities during the post-bounce evolution (see the thin white line in Fig. 2 and the top panel of Fig. 12). At the same time, $z70$ exhibits an accelerated decrease in the minimum lapse function α_{\min} , as illustrated in the bottom panel of Fig. 1, which decreases below $\lesssim 0.01$ at $t_{\text{pb}} \sim 155$ ms and we observe the formation of the apparent horizon. At our final simulation time for $z70$, i.e. less than a millisecond after the apparent horizon appears, its circumferential radius is ~ 2.29 km, with its corresponding mass being $\sim 2.40 M_{\odot}$. Since the initial condition is spherically symmetric, the apparent horizon essentially preserves spherically symmetric property with its deviation from spherical symmetry being less than 0.1%. We note that the size of the quark core at this moment has a radius of ~ 1.8 km, which is well within the apparent horizon and is thus causally disconnected from outside. A comparison with the study of Shibagaki et al. (2021), where the same $z70$ progenitor model was used, shows the black hole formation at a significantly later post bounce time of $t_{\text{pb}} \simeq 230$ ms. This is attributed to the different, purely hadronic EOS of (Lattimer & Swesty 1991) used in the supernova simulations of Shibagaki et al. (2021), which are based on 3D general relativistic neutrino-radiation hydrodynamics simulation. In their case the PNS collapse and black hole formation is induced due to enclosed PNS mass reaching the maximum mass of the hadronic EOS. In contrast here, the PNS collapse is due to reaching the critical conditions for the onset of the phase transition. The phase boundaries, as well as their temperature dependence, are marked as green dashed lines in Fig. 2, for the onset of the hadron-quark mixed phase (left) and for the pure quark matter phase (right). It becomes evident that the evolution proceeds through the hadronic-quark mixed phase into the pure quark matter phase. However, when the second collapse halts in the pure quark matter phase, and even though a second shock wave has formed in compliance with the previous study of Fischer et al. (2018), the enclosed mass of about $2.7 M_{\odot}$ (see Fig. 3) exceeds also here the maximum mass given by the

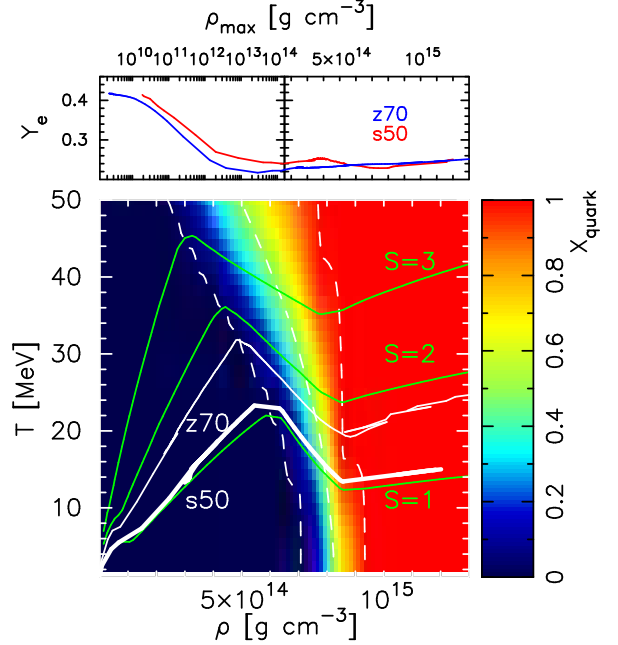


Figure 2. *Top panel:* evolution of the central Y_e as a function of the central rest mass density, ρ_{\max} , for the model $s50$ (red line) and $z70$ (blue line). *Bottom panel:* phase diagram of the DD2F-RDF 1.2 hybrid-EOS, showing the quark volume fraction X_{quark} defined in Bastian (2021), in the density-temperature plane with a fixed electron fraction of $Y_e = 0.25$, which is a typical central value after bounce for both $s50$ and $z70$. The white thick(thin) line denotes the evolutionary path of quantities for $s50(z70)$. The green and white dashed lines are showing the curves of constant entropy per particle of $s = 1, 2, 3 k_B$ and contours of $X_{\text{quark}} = 0.01$ (left), 0.5 (middle) and 0.99 (right), respectively.

DD2F-RDF 1.2 hybrid EOS and a black hole forms eventually.

The $s50$ model takes an intermediate evolutionary path between the less massive ($z9.6$ and $s11.2$) and the most massive ($z70$) progenitor stars. The maximum density increases significantly faster than the lower mass models, due to a higher post-bounce mass accretion rate, though slower than for $z70$. At $t_{\text{pb}} \sim 376$ ms, when the maximum density reaches slightly above two times of the saturation density, the PNS core enters the phase transition region, as illustrated via the thick white line in Fig. 2. Similar as for the $z70$ simulation, also for $s50$ the central PNS becomes gravitationally unstable and collapses once a sufficiently large amount of the PNS interior enters the hadron-quark mixed phase, where the EOS softens substantially in comparison to the pure hadronic and quark matter phases at lower and higher density, respectively. During this second collapse, the maximum density increases rapidly from $\rho_{\max} \simeq 2 \times \rho_{\text{sat}}$ to slightly above $\rho_{\max} \gtrsim 4 \times \rho_{\text{sat}}$ (see Fig. 1). At such a high density, baryons are completely dissociated into quarks and the stiffness of EOS again increases mediated by the repulsive vector interactions. Con-

sequently the NS collapse halts and the second core bounce occurs at $t_{\text{pb}} = 376.6$ ms. The second bounce time seen in the model *s50* is considerably earlier than the finding of about 1.2 s in Fischer et al. (2018). We attribute the reason therefore to the different hybrid EOS applied in Fischer et al. (2018), corresponding to the DD2–RDF 1.1 nomenclature of Bastian (2021). It features a generally higher onset density for the hadron-quark phase transition, on the order of $3.5 \times \rho_{\text{sat}}$ at $T = 0$, as well as a different temperature dependence. Further differences in the simulations may arise due to the different hydrodynamics scheme of the supernova model of Fischer et al. (2018), which is based on an adaptive Lagrangian mass grid and differs from the one in this study. Furthermore, we have checked if the spatial resolution can influence on the evolution up to and including second bounce and found a rough numerical convergence, which is discussed in Appendix A.

Figure 2 depicts how the central PNS evolves in the phase diagram of the DD2F–REF 1.2 hybrid EOS. The top panel shows the evolution of the central Y_e as a function of the central rest mass density for the models *s50* (red line) and *z70* (blue), with $Y_e \simeq 0.25$ for both simulations. The bottom panel exhibits the phase diagram of hybrid EOS employed with the color-coding indicating the quark volume fraction, X_{quark} , in the density-temperature plane for a fixed $Y_e = 0.25$. The white thick and thin solid lines denote evolutionary paths of the central density and temperature for *s50* and *z70*, respectively. The green lines show curves of constant entropy per particle of $s = 1, 2, 3 k_B$, and the vertical white dashed lines correspond to $X_{\text{quark}} = 0.01, 0.5$ and 0.99 .

In both models *s50* and *z70*, the central region evolves nearly along a constant entropy path after the first bounce due to the adiabatic compression. It lasts until it reaches the purely quark matter phase, i.e., where $X_{\text{quark}} = 1$. From the phase diagram, we clearly see a temperature dependence of the initiation of the hadron quark mixed phase, i.e. the left vertical white dashed curve. It shows that the mixed phase appears at lower density for matter with higher temperature or entropy. This turns out that the model *z70*, possessing higher core entropy, enters the phase transition region at a slightly lower density of $\rho_{\text{max}} \simeq 5 \times 10^{14} \text{ g cm}^{-3}$, compared to the value of $\rho_{\text{max}} \simeq 6 \times 10^{14} \text{ g cm}^{-3}$ in *s50*. After the central region enters the pure quark matter phase, we witness multiple core bounces for *z70*, due to the stiffening of EOS that produce a standing shock front. However, the mass accretion rate is extremely high, on the order of ($\sim 4 M_{\odot} \text{ s}^{-1}$, measured on the shock surface at the second collapse, and the shock front cannot propagate outward to blow off the outer envelop. Instead ρ_{max} increases continuously and a black hole is quickly formed as already explained. On the other hand, the maximum density for *s50* does not continuously increase once its PNS interior transforms into a quark matter core. In

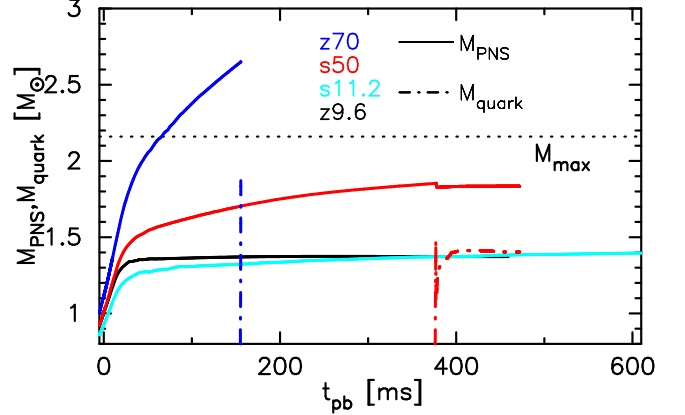


Figure 3. Post-bounce evolution of the rest mass of the PNS mass M_{PNS} (solid lines) and proto-quark star mass M_{quark} (dashed) for all models. The horizontal dotted line marks the maximum mass at $T = 0$ of $M_{\text{max}} = 2.16 M_{\odot}$ for the current EOS. M_{quark} for *z9.6* and *s11.2* is zero since the hadron-quark phase transition region is not reached during our simulation time of $t_{\text{pb}} \lesssim 600$ ms.

Fig. 1, ρ_{max} shows a single bounce and then quickly settles into a new value of $\rho_{\text{max}} \simeq 1.05 \times 10^{15} \text{ g cm}^{-3}$ after a rapid ring down phase. At the same time, the minimum lapse α_{min} also decreases from ~ 0.7 to ~ 0.5 across the second core bounce. The second bounce creates an energetic hydrodynamic shock wave due to the stiffening of EOS. In contrast to *z70*, which has a significantly higher mass accretion rate, the shock front propagates outward with a relativistic velocity of about two-thirds of the speed of light, because of a lower mass accretion rate of about $0.33 M_{\odot} \text{ s}^{-1}$ measured on the shock surface at the second core collapse. The shock evolution will be discussed in the next section 3.2.

Qualitative difference between the BH formation model (*z70*) and the explosion model (*s50*) can be explained from evolution of PNS mass. In Fig. 3, we plot the PNS mass M_{PNS} (solid lines) and proto-quark star mass M_{quark} (dashed) for all models. The latter is defined via the quark volume fraction exceeding a few percent. It has been introduced in Fischer et al. (2018) and further details, including the phase transition construction, can be found in Bastian (2021). We define the PNS surface by the region where the rest mass density drops below $10^{10} \text{ g cm}^{-3}$.

In contrast to the *z70* model, the baryon mass enclosed inside the PNS for *s50* of about $1.83 M_{\odot}$ is well below the maximum gravitational mass of DD2F–RDF 1.2 after the hadron-quark phase transition (see Fig. 3). Note that the maximum allowed mass for hot neutron stars can be increased by $\sim 0.2\text{--}0.4 M_{\odot}$, compared to their corresponding cold case due to the corresponding additional pressure support from hot matter (c.f. Hempel et al. 2012). The nascent PNS for *s50* features a massive quark matter core of about $1.43 M_{\odot}$. Similar was the case for *z70*, a strong second hydrodynamics

shock wave forms as a result of the phase transition. However, in the case of *s50* the remnant PNS is semistable, e.g., with respect to radial and non-radial perturbations such that a black hole formation is not observed. The dynamics of the second shock wave has important consequences for the supernova explosion dynamics, which will be discussed below.

None of the lighter progenitor models, *z9.6* and *s11.2*, reach the conditions required for the hadron-quark phase transition in our simulation times of less than ~ 600 ms after bounce. Instead, these two models show the explosion driven by neutrino heating. This will be explained in details in the next sec. 3.2. It is noteworthy that the maximum density continues increasing slightly after the explosion begins. If the central density increases with the same rate at the final simulation time $\dot{\rho}_{\max} \sim 10^{14} \text{ g cm}^{-3} \text{ s}^{-1}$, we estimate that it will exceed $2 \times \rho_{\text{sat}}$ at $t_{\text{pb}} \sim 1.5$ s in both models and the hadron-quark phase transition may occur during the subsequent deleptonization phase of the nascent NS. However, we note here also that the phase transition onset mass of DD2F-RDF 1.2 for cold neutron stars, i.e. the mass threshold for the hybrid star branch lies above $1.37 M_{\odot}$. Hence, in case the phase transition would occur we expect that only a tiny fraction, if any, of the neutron star’s interior would consist of pure quark matter.

In Table 1 we list a summary of progenitor models and their final fates. The compactness parameter ξ_M measures the ratio of mass $M(M_{\odot})$ and radius R_M , where the latter term R_M expresses the radius, which encloses the mass M (O’Connor & Ott 2011), and is in units of 1000 km. Regarding the compactness of progenitor star, ξ , *z70* shows the highest value. O’Connor & Ott (2011) discussed a relation between the bounce compactness $\xi_{2.5}^{\text{b}}$ and BH formation time. According to their result, BH formation time is less than a few 100 ms for very massive stars having $\xi_{2.5}^{\text{b}} \gtrsim 1.5$, which is in accordance with the current model *z70*. *s50* forms a hybrid star with its quark core and total PNS mass reach 1.43 and $1.83 M_{\odot}$, respectively. *z9.6* and *s11.2*, which have a small compactness parameter, show almost the same PNS mass $1.37 M_{\odot}$ with their evolution showing a plateau behavior (see black and cyan lines in Fig. 3).

3.2. Explosion dynamics

In this section we discuss the explosion dynamics for models *z9.6*, *s11.2*, and *s50*. As *z70* does not explode, its E_{exp} is zero. Fig. 4 exhibits time evolution of averaged shock radii R_{shock} (top panel) and diagnostic explosion energy E_{exp} (bottom). We note that E_{exp} does not take into account the binding energy of the stellar envelop. In addition nuclear recombination and the associated nuclear burning energy generation are also neglected.

We begin by *z9.6*. From its shock evolution (black curve in Fig. 4), we see that the shock expands slowly at initial

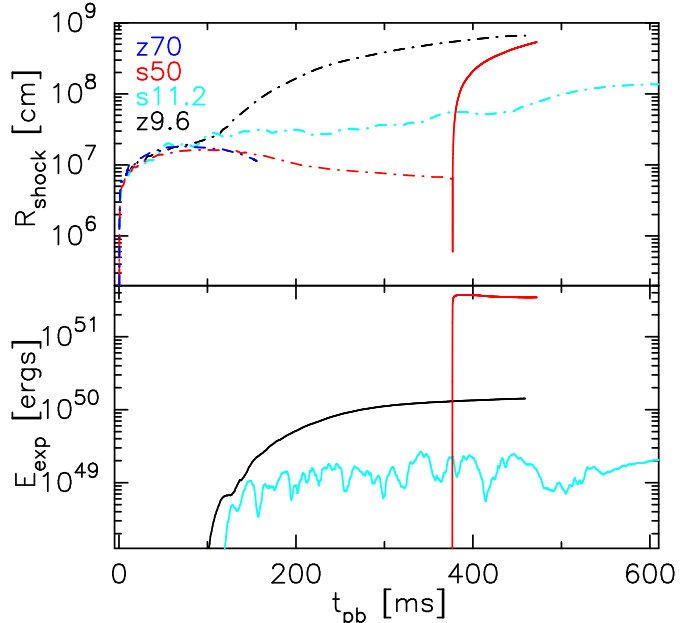


Figure 4. Post-bounce evolution of the averaged shock radii R_{shock} (top panel) and of the diagnostic explosion energy E_{expl} (bottom panel) for all models. In the upper panel, we distinguish the core bounce shock (denoted by the dash-dotted line) and the second bounce shock (solid).

times up to about 100 ms post bounce. After about 120 ms post bounce, this model enters a shock runaway phase. The shock front reaches 1000 km at $t_{\text{pb}} = 167$ ms. At the same time, the diagnostic explosion energy E_{exp} starts increasing and eventually reaches a value of 1.4×10^{50} erg when the shock front reaches the outer boundary of the calculation domain. In Melson et al. (2015b) and Radice et al. (2017), the authors also performed multi-D SN simulations using the same $9.6 M_{\odot}$ progenitor star. Although the detailed numerical setups are different in each study, e.g., different hydrodynamics schemes, grids, neutrino transport and weak rates, the shock and diagnostic explosion energy evolution show quantitative behavior. Both Melson et al. (2015b) and Radice et al. (2017) report $E_{\text{exp}} \sim 8 \times 10^{49}$ erg at $t_{\text{pb}} = 400$ ms, where the latter study takes into account the gravitational binding energy of the envelope, which is merely -2×10^{48} erg (c.f. Melson et al. 2015b). We thus find consistent results with previous studies for model *z9.6*.

For the other low mass model *s11.2*, the average shock radius steadily increases without any visible shock stagnation, similarly to *z9.6*, however, at a somewhat slower rate. This is attributed to the extended silicon-sulphur layer for this progenitor model, which is absent for *z9.6*. It results in higher mass accretion rates post bounce for *s11.2*, which, in turn slow down the neutrino driven shock expansion. Fig. 4 shows that the mean shock radius R_{shock} reaches 1000 km at

model	$\xi_{1.5}^{b(i)}$	$\xi_{2.5}^{b(i)}$	M_{remnant} [M_{\odot}]	M_{quark} [M_{\odot}]	t_{exp} [ms]	t_{final} [ms]
z9.6	(2.2×10^{-4})	(7.6×10^{-5})	1.37	—	167	459
s11.2	$2.4(1.8) \times 10^{-1}$	(5.0×10^{-3})	1.40	—	503	650
s50	2.3(1.6)	$1.9(1.9) \times 10^{-1}$	1.83	1.35	384	450
z70	4.0(1.1)	1.6(1.0)	2.65	1.88	—	155

Table 1. List of models and their simulation results, showing the compactness parameters, ξ_M , measured at two different mass coordinates of $1.5 M_{\odot}$ and $2.5 M_{\odot}$, the remnant masses, M_{remnant} , the mass of the quark-matter core, M_{quark} , the explosion time, t_{exp} , and the final simulation time, t_{final} . We evaluate ξ_M at core bounce and at the initial data, denoted by ξ_M^b and ξ_M^i (values in parantheses), respectively. Note that our calculation domain does not contain $1.5 M_{\odot}$ for z9.6 and $2.5 M_{\odot}$ for s11.2 and we omit the corresponding compactness parameters at bounce ξ_M^b from the table. Explosion time t_{exp} denotes the post bounce time when the average shock radius reaches $r = 1000$ km.

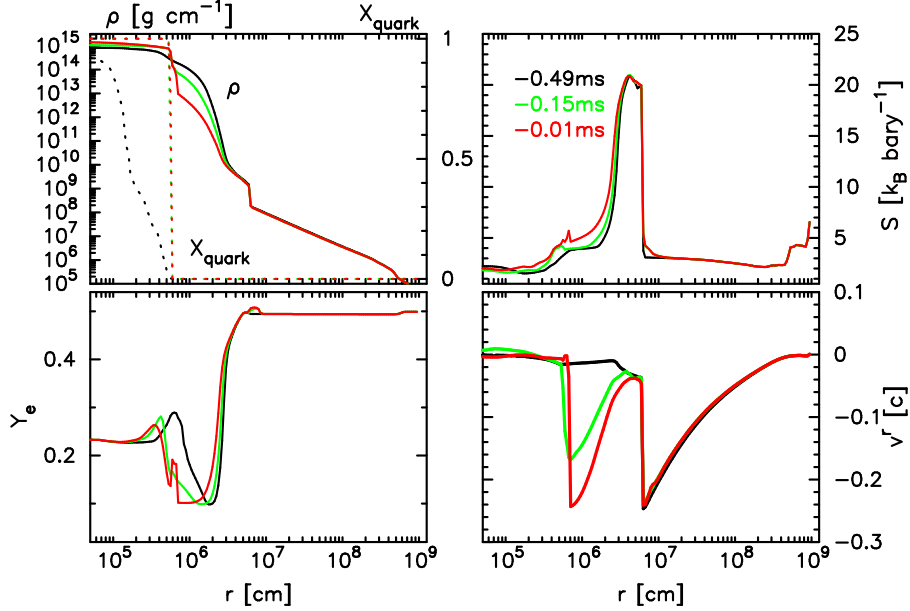


Figure 5. Radial profiles of the rest mass density ρ (solid lines) and quark fraction X_{quark} (dotted), entropy per particle s , electron fraction Y_e , and radial component of the three-velocity $v^r \equiv u^r/u^t$ for the model s50 at different times before to the second bounce, denoted as t_{pb} , corresponding to the onset and proceeding collapse of the PNS due to the hadron-quark phase transition.

$t_{\text{pb}} \sim 503$ ms and the diagnostic explosion energy reaches $3\text{--}4 \times 10^{49}$ erg at the end of simulation. These values are in quantitative agreement with those in Müller (2015), in which they conducted a relativistic 2D CCSN simulation for the same $11.2 M_{\odot}$ progenitor model. Although the employed EOS and neutrino opacities are different, and comparison is not straightforward, his corresponding 2D model shows a similar evolution with s11.2 in this study. The neutrino driven supernova explosion observed for this low mass model s11.2 is in qualitative agreement with the results reported by other groups based on axially symmetric simulations (c.f. Buras et al. 2006; Marek & Janka 2009; Takiwaki et al. 2014; Nagakura et al. 2018). More controversial are the recent 3D supernova simulations of this progenitor model (e.g., Takiwaki et al. 2012; Tamborra et al. 2014; Takiwaki et al. 2016), which is an active subject of research.

In comparison to these less massive progenitor stars, s50 shows a continuous shock recession once the shock stagnates around $t_{\text{pb}} \sim 100$ ms. The recession lasts until the second collapse occurs (see Fig. 4). The red line in the top panel shows that the averaged shock radius eventually shrinks to $R_{\text{shock}} = 62$ km. Then concomitantly with the second bounce, a strong, second shock wave is produced at the surface of quark core, locating at around $r \sim 7$ km. This shock front overtakes the preceding standing shock front at $r = 62$ km and propagates further out with a relativistic velocities, i.e. within a few ms after the second bounce, the shock front exceeds $r = 1000$ km. Simultaneously the diagnostic explosion energy rises continuously, and reaches its asymptotic value of $E_{\text{exp}} \sim 3.5 \times 10^{51}$ erg. Afterward, E_{exp} shows nearly a constant value throughout the runaway explosion phase. We estimate that the gravitational binding energy of the envelop is roughly -4.0×10^{50} erg and thus the asymptotic explo-

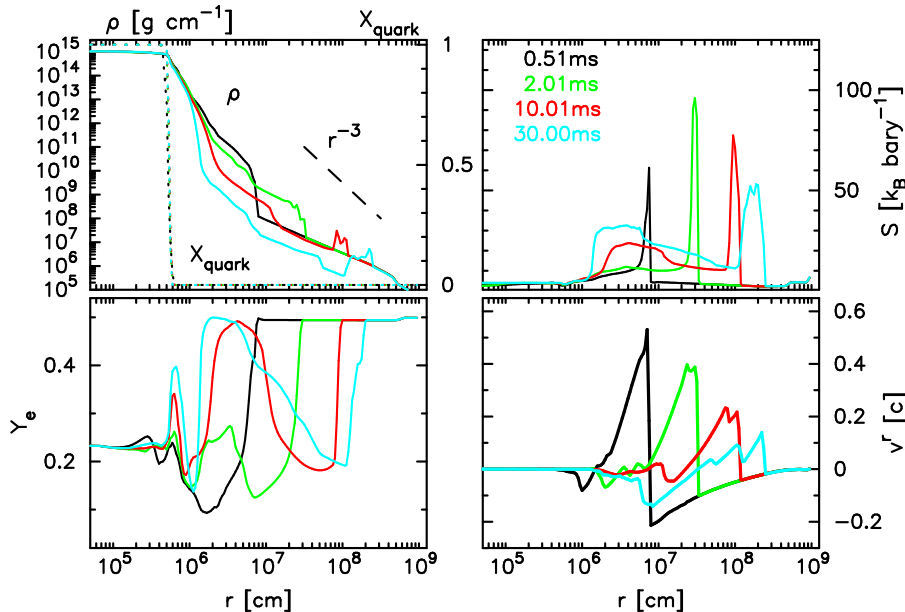


Figure 6. The same as Fig. 5, but for different time slices after the second bounce. In the density profile, we show a slope of r^{-3} by a dashed line.

sion energy would be around 3×10^{51} erg. In view of the explosion dynamics, similar as was discussed in Fischer et al. (2018), our model *s50* supports that these massive progenitor stars could be the progenitors of a kinds of superluminous supernova events, (c.f. the discussion in Smith et al. 2007; Moriya et al. 2013, and references therein), if a sufficiently strong hadron-quark phase transition occurs.

Next, to illustrate the dynamics at the second collapse and in the subsequent explosion phase, we plot spherically averaged profiles of the rest mass density ρ and quark fraction X_{quark} (top left panel), entropy s (top-right panel), electron fraction Y_e (bottom left panel), and the radial component of the three-velocity v^r (bottom right panel) in Figures 5 and 6, at selected time slices.

Fig. 5 shows profiles before second bounce $t_{\text{p2b}} < 0$, where t_{p2b} denotes the post second bounce time and written in the top left panel. From the density and radial velocity profiles, we see that the second collapse takes place within a fraction of a millisecond. At $t_{\text{p2b}} = -0.49$ ms (black lines), the radial velocity profile shows nearly a hydrostatic profile ($v_r \sim 0$) with a slight infalling sign inside the PNS. At the same time, quarks are already partially deconfined, showing $X_{\text{quark}} \sim 0.9$ at the center. We note that quarks are completely confined in hadronic-matter, i.e. $X_{\text{quark}} = 0$, up until ~ 1 ms before the second bounce, which is a major difference from Zha et al. (2020). Shortly after that, the PNS core collapses with a relativistic speed of $\sim 0.25c$ (see the red line in v^r profile) and the PNS core transforms into a quark core with $X_{\text{quark}} = 1$. While the entropy (top-right) and Y_e (bottom-left) do not show any noticeable changes compared to X_{quark} and v^r .

After the second bounce happens, the hydrodynamic profiles change considerably. The radial velocity profile shows a relativistic outward shock propagation immediately at $t_{\text{p2b}} = 0.51$ ms and, already at this moment, the entropy behind the shock reaches $s \geq 60 k_B$ baryon $^{-1}$. In the subsequent phase, however, its propagation speed decreases gradually as clearly be seen in the bottom right panel. At $t_{\text{p2b}} = 30$ ms, the shock decelerates to $\sim 0.1c$ (cyan line in the bottom-right panel). The entropy behind the shock also decreases noticeably from $\sim 100 k_B$ baryon $^{-1}$ at $t_{\text{p2b}} = 2$ ms to $\sim 50 k_B$ baryon $^{-1}$ at $t_{\text{p2b}} = 30$ ms. Furthermore a density bump is formed around $r \sim 10^8$ cm at $t_{\text{p2b}} = 10$ ms. All of these facts indicate that the shock front propagates in a dense low-entropy medium and is substantially decelerated. Indeed from the top left panel in Fig. 6, the angle averaged density profile shows a shallower density gradient than the slope of r^{-3} at 10^7 cm $\lesssim r \lesssim 5 \times 10^8$ cm. In such a dense environment, the shock can be decelerated very efficiently (Sedov 1959; Wongwathanarat et al. 2015), resulting in development of Rayleigh-Taylor (RT) instabilities, which we will describe later.

The Y_e profile also shows a remarkable feature. The low- Y_e materials, which initially locate in the Y_e -trough just above the PNS surface $r \sim 20$ km (see black line in the bottom-left panel), are ejected along with the explosion, since the explosion initiates on the quark core surface at $r \sim 7$ km, i.e. underneath the Y_e -trough. At $t_{\text{p2b}} = 30$ ms, those low- Y_e ejecta are reaching $r \sim 10^8$ cm and produce a steep Y_e composition gradient across the shock surface. We also find a high- Y_e and high entropy region, for instance, with $s \sim 30 k_B$ baryon $^{-1}$ and $Y_e \sim 0.5$ around $r \sim 5 \times 10^6$ cm at $t_{\text{p2b}} \geq$

30 ms. These components are produced due to a powerful neutrino irradiation.

A combination of the aforementioned shock deceleration and steep gradient of Y_e results in a very remarkable multi-D hydrodynamic feature. To illustrate it, we depict 2D profiles of the shocked region at several time slices in Fig. 7. We select four time slices: $t_{\text{pb}} \simeq 150$ ms as a representative of normal post-bounce structure in Fig. 7(a) and $t_{\text{pb(p2b)}} = 390(14)$ in Fig. 7(b), 416(40) in Fig. 7(c), and 461(85) ms in Fig. 7(d) as example after the second bounce. Each panel consists of four panels, in which we plot: logarithmic density $\log_{10}(\rho [\text{g cm}^{-3}])$ (in top-left mini panel), entropy per particle $s [k_{\text{B}}]$ (top-right), temperature $T [\text{MeV}]$ (bottom-left), and electron fraction Y_e (bottom-right) for model *s50*.

Top-left panel shows normal convective motions driven mainly by neutrino heating. At $t_{\text{pb}} = 150$ ms, the shock front locates $r \sim 140$ km and high entropy ($s \sim 15 k_{\text{B}}$ baryon $^{-1}$) convective bubbles are seen behind it. In the post second bounce phase, the shock expands initially with a relativistic speed, though it starts decelerating already at $t_{\text{p2b}} \gtrsim 10$ ms. Indeed from the density contour at $t_{\text{pb(p2b)}} = 390(14)$ ms (top-right panel), we see a dense shell locating between $10^8 \text{ cm} \lesssim r \lesssim 1.5 \times 10^8 \text{ cm}$. Within the shell, its rest mass density exceeds $\rho = 10^7 \text{ g cm}^{-3}$, which is approximately one order of magnitude larger than the value of its surrounding. Particularly interesting feature is seen on the inner edge of the shell at $r \sim 10^8 \text{ cm}$, where the density has a negative radial gradient. The deceleration of the dense shell produces an effective outward acceleration in the comoving frame, acting on the shell component. Therefore it fulfills the condition for development of the Rayleigh-Taylor instability. As a consequence, higher density materials, which initially locate more inward, penetrate into the lower density side and the well known mushroom like structures are formed. Those mushroom like plumes are composed of neutron rich material with $Y_e \lesssim 0.3$ (see bottom-right minipanel at $t_{\text{pb(p2b)}} = 390(14)$ ms).

From Figs. 7(c) and 7(d), one can find how effectively those plumes, associated with the Rayleigh-Taylor instability, transport the low- Y_e material outward. The low- Y_e plumes are growing at $r \gtrsim 10^8 \text{ cm}$. At such a distant region from the central hot PNS, neutrino irradiation is significantly weak and neutrino absorption basically does not occur. Consequently Y_e of the ejecta does not basically change except through the hydrodynamic mixing process. The radial three-velocity of the plumes is roughly $v^r \sim 0.1c$ at our final simulation time of $t_{\text{pb(p2b)}} = 475(100)$ ms and approaches nearly a steady state value. Fischer et al. (2020a) reported *r*-process nucleosynthesis calculation based on spherically symmetric explosion models triggered by the PDE mechanism. According to their results, the explosion can indeed transport low- Y_e material to the distant region $r \gtrsim 10^{8-9} \text{ cm}$, though the Y_e value tends to

become higher $\gtrsim 0.35$ compared to this study. This is simply due to the absence of the RT mixing under the spherically symmetry. Therefore, with the aid of the RT mixing, the low- Y_e components can be more efficiently ejected and may significantly influence on the nucleosynthetic yields, which will be reported elsewhere.

The Rayleigh-Taylor instability can be seeded not only by the inhomogeneous structure behind the standing shock surface (see the top-left panel at $t_{\text{pb}} \sim 150$ ms in Fig. 7), but also by the convective motions inside the quark core itself. To explain this, we evaluate the Brunt-Väisälä frequency ω_{BV} and plot its spherically averaged spatial profile around the second bounce in Fig. 8. Here, we follow the same manner as described in Müller et al. (2013) to evaluate ω_{BV} (see their Eq. (C11)) in relativistic form. In the figure, we select two time slices across the second bounce, represented by red ($t_{\text{p2b}} = -0.01$ ms) and blue ($t_{\text{p2b}} = 0.05$ ms) lines. In these time snapshots, the quark core surface, where $X_{\text{quark}} = 1$, locates at $r \sim 6 \times 10^5 \text{ cm}$, without changing its location. While the shock front propagates from $r \sim 6 \times 10^5 \text{ cm}$ at $t_{\text{p2b}} = -0.01$ ms to $\sim 9 \times 10^5 \text{ cm}$ at $t_{\text{p2b}} = 0.05$ ms. It is obvious that ω_{BV} in front of the second bounce shock shows positive, indicating convectively unstable region from the Ledoux criterion, and high angular frequencies, reaching $\sim 10 \text{ ms}^{-1}$. The convectively unstable region corresponds to the region with a steep negative lepton number gradient (see Y_e profile in Fig. 5) as well as with the entropy jump at the second bounce shock surface. Due to such a high ω_{BV} , the convective motions can fully develop behind the shock even within a millisecond and be the seed of the subsequent Rayleigh-Taylor instability. Furthermore, in the convectively stable region (i.e. $\omega_{\text{BV}} < 0$) between the quark core and the second bounce shock surface ($6 \times 10^5 \text{ cm} \lesssim r \lesssim 9 \times 10^5 \text{ cm}$ at $t_{\text{p2b}} = 0.05$ ms), we see the angular frequencies reaching $\omega_{\text{BV}} \sim$ a few times -1 ms^{-1} . This frequency $|\omega_{\text{BV}}|$ = a few kHz, i.e. the frequency of convection motions, can reasonably explain the GW frequencies at the second bounce, which will be discussed in the next section.

4. GW AND NEUTRINO EMISSION

In this section we discuss the impact of the hadron-quark phase transition on the emission of gravitational waves and neutrinos. In Fig. 9, we plot gravitational waveforms and their spectrogram for all models. Here, we show only a non-vanishing component $A_+ = Dh_+$ observed along the equatorial plane that is perpendicular to the symmetric axis. D and h_+ are the source distance and gravitational wave strain, respectively. h_+ is evaluated from a standard quadrupole formula (Shibata & Sekiguchi 2003; Kuroda et al. 2014).

A common feature seen among all models is relatively large GW amplitudes, observed after the first core bounce till $t_{\text{pb}} \sim 50$ ms. These GWs are emitted from the postbounce

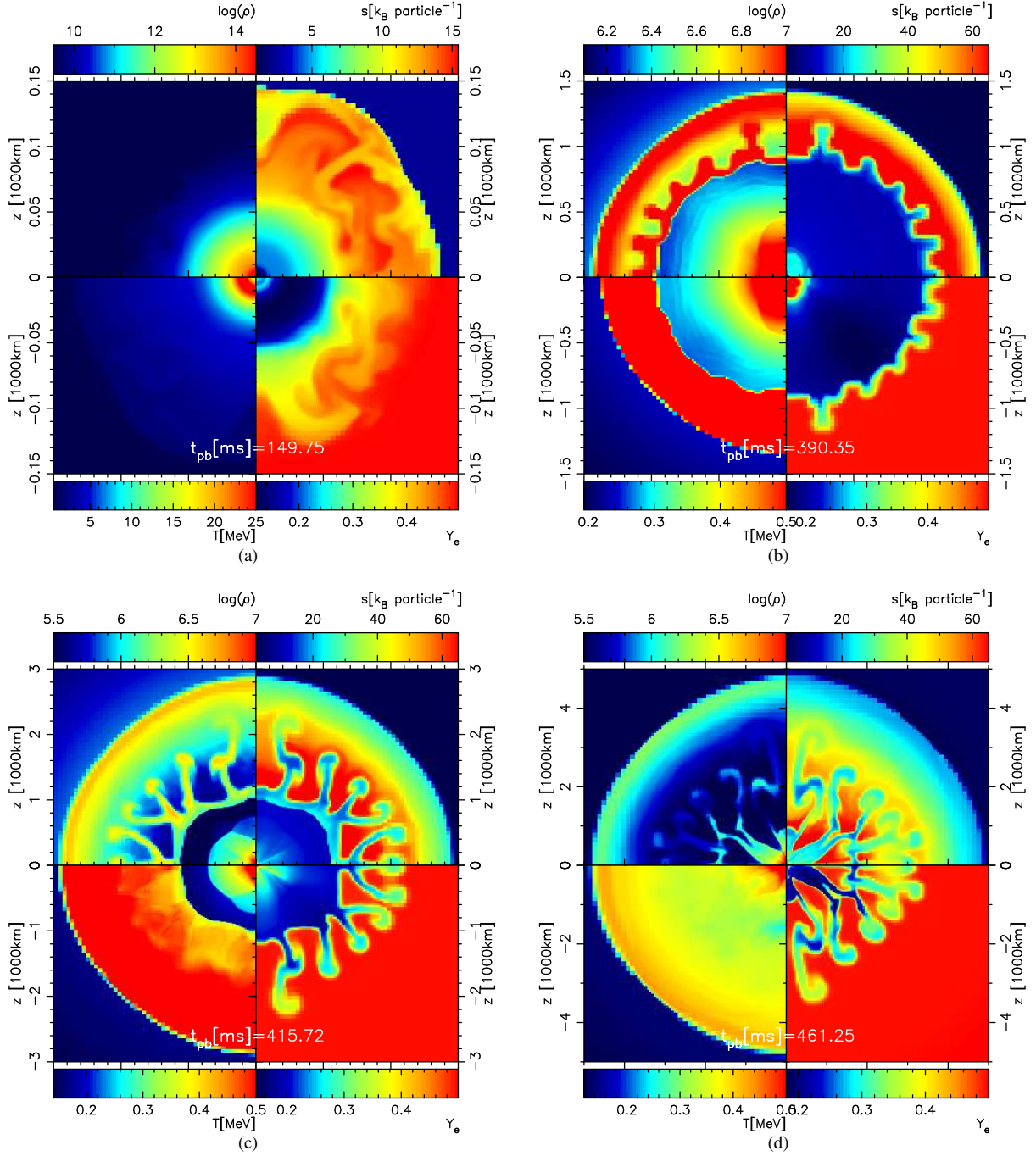


Figure 7. Snapshots of the $x - z$ -plane for the logarithmic density $\log_{10}(\rho)$ [g cm^{-3}] (top-left panels), entropy per particle s [k_B] (top-right panels), temperature T [MeV] (bottom-left panels), and Y_e (bottom-right panel) for $s50$, at different time slices $t_{pb} = 150$ ms in graph (a) as a representative of the normal structure during the post bounce phase prior to the hadron-quark phase transition as well as at $t_{pb}(t_{p2b}) = 390(14)$ in graph (b), 416(40) in graph (c) and 461(85) ms in graph (d) after the second bounce.

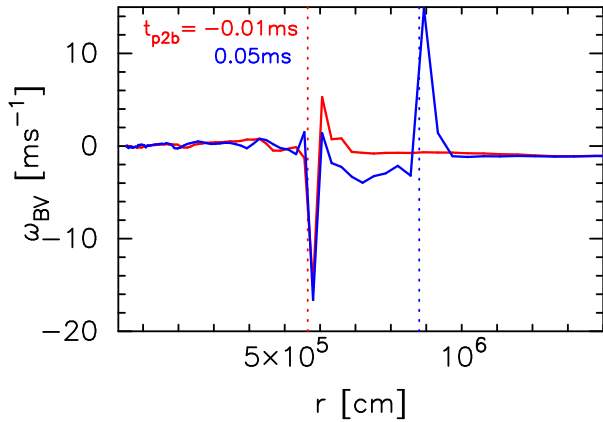


Figure 8. Spherically averaged spatial profiles of the Brunt-Väisälä frequency ω_{BV} for $s50$ around the second bounce. Here we focus only on the region above of the quark-matter core and we select two representative time slices at $t_{\text{p2b}} = -0.01$ ms (red line) and $t_{\text{p2b}} = 0.05$ ms (blue line). The vertical dashed lines indicate the averaged radii of the second bounce shock. Positive ω_{BV} indicates convectively unstable region based on the Ledoux criterion.

convective motions and their amplitudes reach $\sim 20 - 50$ cm. Afterward, the gravitational waveforms show various features depending on the progenitor mass. Our least massive star $z9.6$, which is the NDE model and experiences the runaway explosion at around $t_{\text{pb}} \sim 100$ ms (see, Fig. 4), shows a considerable subsidence of GWs. The waveform is quite similar with that of a 2D relativistic model of (Müller et al. 2013).

Another NDE model $s11.2$ presents continuous GW emission even after the shock front exceeds 10^8 cm at $t_{\text{pb}} \sim 500$ ms. Furthermore, the waveform seems to show an increasing frequency, which can be indeed confirmed from the spectrogram in the lower panel. The peak frequency concentrates around ~ 200 Hz at $t_{\text{pb}} = 200$ ms and is increasing with time. Such a ramp-up feature of the GW frequency is widely seen in CCSN models (Murphy et al. 2009; Müller et al. 2013; Kuroda et al. 2016a; Vartanyan et al. 2019; Mezzacappa et al. 2020; Shibagaki et al. 2021) and may come from the g/f -mode oscillation of the non-rotating PNS (Morozova et al. 2018; Torres-Forné et al. 2019; Sotani & Takiwaki 2020) or perhaps PNS convection (Mezzacappa et al. 2020). In the late post bounce phase $t_{\text{pb}} \gtrsim 600$ ms, the spectrogram shows a wide band emission ranging from 300 Hz to a few kHz. Since we use a simple short time Fourier analysis to obtain the spectrogram, we need to apply sophisticated analysis methods, such as S-method (Kawahara et al. 2018), the Hilbert-Huang transform (Takeda et al. 2021), or non-harmonic analysis (Yanagisawa et al. 2019) to determine more detailed peak frequency.

Regarding the PDE model $s50$, the ramp-up feature can also be seen in the first ~ 200 ms after bounce. It is remarkable

that the ramp-up mode is showing a more rapid evolution in its peak frequency than that in $s11.2$. At $t_{\text{pb}} \sim 180$ ms, the peak frequency appears at ~ 500 Hz, which is significantly higher than ~ 200 Hz for $s11.2$. The higher frequency in model $s50$ is simply due to a higher compactness of the PNS (see Table 1) and its higher g/f -mode oscillations. From $t_{\text{pb}} \sim 200$ ms to the second bounce time $t_{\text{pb}} \sim 380$ ms, the GW emission shows a quiescent phase, during which the amplitudes are basically less than ~ 1 cm. Such a quiescent phase can also be seen in our most massive star case $z70$ at $t_{\text{pb}} \gtrsim 100$ ms. In both models, the quiescent phase roughly corresponds to the period when the shock gradually recedes (Fig. 4). Although this paper does not aim to clarify the physical mechanism of the quiescent phase, we speculate that the rapid shock recession may suppress the convective motions, which usually excite the PNS core g/f -mode oscillations. As a consequence, the weaker GWs are observed. However, in the full 3D models, the situation may drastically change, as another degree of spatial freedom allows for the development of the standing accretion shock instability (SASI) (Blondin et al. 2003). In such models, the more compact PNS core is more susceptible to the SASI development. The SASI activity can then imprint its feature in the emitted GW with sizeable amplitudes, typically in the low frequency range of $\sim 100 - 200$ Hz (Kuroda et al. 2016a; Andresen et al. 2017; Vartanyan et al. 2019; Mezzacappa et al. 2020; Shibagaki et al. 2021).

The quiescent phase is disrupted by the hadron-quark phase transition. Both models $s50$ and $z70$ experience the second collapse at $t_{\text{pb}} \sim 380$ ms and $t_{\text{pb}} \sim 155$ ms, respectively. Although $z70$ immediately forms a BH, both models produce a quark core at their center and exhibit core bounce and subsequent ring down phase. In inset of the upper panel for $s50$, we show a magnified view of the waveform at the second bounce. The x -axis measures the time relative to the second bounce in millisecond. The GW amplitudes reach $A_+ \sim 250$ cm and the strong emission lasts about 4 ms after the second bounce time. Its burst like feature results in a wide band emission, ranging from ~ 500 Hz to beyond 2.5 kHz, where the latter value corresponds to the Nyquist frequency of our sampling rate of 5 kHz. Such high frequencies are comparable with typical frequencies of the convection motions behind the second bounce shock as shown in Fig. 8, where we see $|\omega_{\text{BV}}| \sim$ a few kHz at $6 \times 10^5 \text{ cm} \lesssim r \lesssim 9 \times 10^5 \text{ cm}$. Therefore these high-frequency GWs could be an outcome of these convection motions induced by the hadron-quark phase transition. During the post second bounce phase in $s50$, the GWs show a high frequency emission at the range of $\gtrsim 1$ kHz and they subside after $t_{\text{pb(p2b)}} \sim 460(80)$ ms. The maximum amplitude and the duration time of the strong emission are comparable to those in Zha et al. (2020).

Regarding the BH formation model $z70$, it also shows a high peak frequency $\gtrsim 1$ kHz and their amplitudes reach

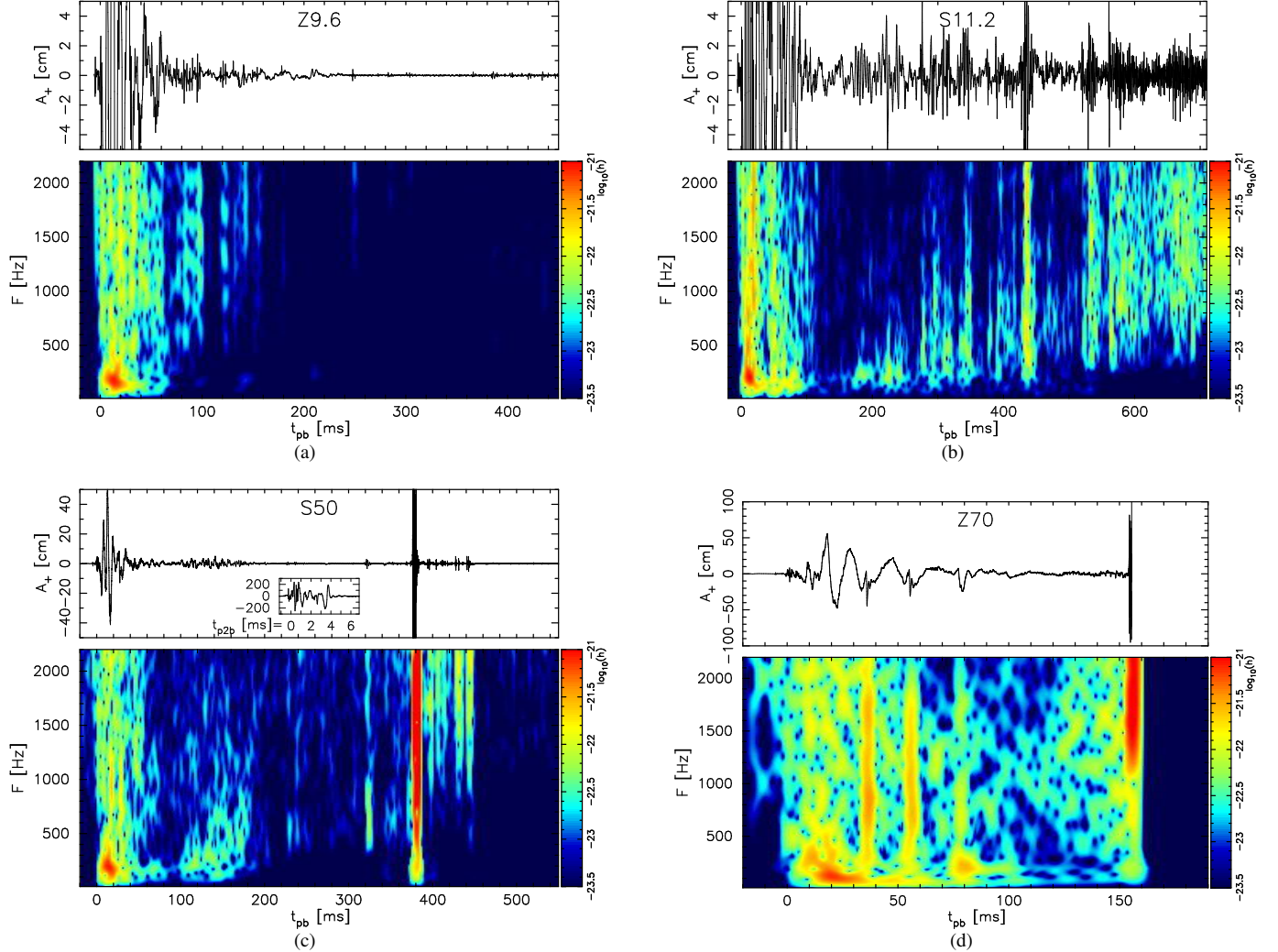


Figure 9. Gravitational waveform A_+ (top panels) and the corresponding spectrogram (bottom panels) for all models, labeled in each top-panel. Here, we show only the non-vanishing component in axisymmetric profile $A_+ \equiv Dh_+$ observed along the equatorial plane, which is perpendicular to the symmetric axis. The spectrogram is obtained by short-time Fourier transform. D and h_+ are the source distance and gravitational wave strain, respectively, where the latter is calculated from a standard quadrupole formula. For *s50*, we also plot a magnified view of the gravitational waveform with respect to the second bounce time in the inner mini-panel of the top-panel.

$A_+ \sim 100$ cm. In contrast to the full 3D BH formation models (Kuroda et al. 2018; Shibagaki et al. 2021), whose GW waveforms are contaminated by the low-frequency SASI modes, the GW signal at the BH formation shows a higher peak frequency (this could be consistent with Sotani & Sumiyoshi 2019, 2021). This is not only due to the absence of significant SASI motions but also to the quark core formation and its oscillation. We crudely check if the GW emission prior to the BH formation actually emanate from outside the apparent horizon. In Fig. 10, we show the GW amplitudes A_+ for model *z70* contributed from different spherical shells: innermost region with $r \leq 3$ km (red line), $3 \text{ km} < r \leq 20$ km (green), and $20 \text{ km} < r \leq 300$ km (blue). Black line indicates the total gravitational waveform. To plot the figure, we simply divide the computational domain into several spheri-

cal shells, calculate the contained mass quadrupole moment inside each shell, and derive the gravitational wave strain.

One can see that the blue line overlaps with the black one in most of the post bounce phase up to $t_{pb} = 150$ ms. It indicates that the dominant contribution to the total GW is coming from the post shocked region, roughly corresponding to the region with $20 \text{ km} < r \leq 300$ km. While the inner region with $r \leq 20$ km shows essentially no GW emission except during the initial post convection phase ($t_{pb} \lesssim 30$ ms). However, just prior to the second collapse ($t_{pb} \geq 154$ ms, i.e. less than ~ 1 ms till the BH formation), the dominant contribution is from $r \leq 20$ km. At this moment, the quark core is already formed. Although the contribution from the innermost region with $r \leq 3$ km is not negligible (see the red line), the majority comes from outside of the apparent horizon, i.e. 3 km

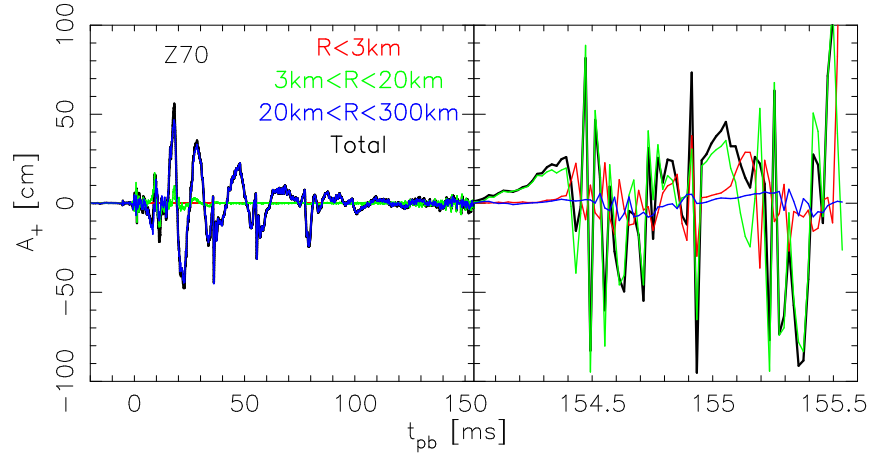


Figure 10. GW amplitudes A_+ for the model $z70$ contributed from different spherical shells: innermost region with $r \leq 3$ km (red line), $3 \text{ km} < r \leq 20$ km (green line), and $20 \text{ km} < r \leq 300$ km (blue line). Black line indicates the total gravitational waveform.

$< r \leq 20$ km. Here we note that the circumferential radius of AH is ~ 2.29 km at the final simulation time. Therefore, the signal of the QCD phase transition could be imprinted into the emanated GWs, even in the failed CCSN models. We, however, stress that the gravitational waveform can greatly change in the full 3D models and also that the gravitational wave should be evaluated more precisely by the gauge invariant method, especially in such a model on a highly curved spacetime.

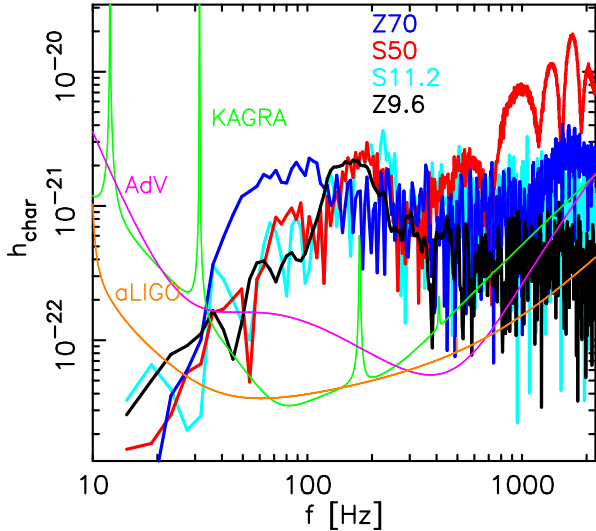


Figure 11. Characteristic GW spectral amplitudes for all models assuming a source distance of 10 kpc. The noise amplitudes of aLIGO (orange line), AdV (magenta line), and KAGRA (green line) are plotted as reference. Note that, since the Nyquist frequency is 2500 Hz, we show only $f \lesssim 2000$ Hz.

To discuss the detectability of the GW signals, we show the GW spectral amplitudes for all models in Fig. 11, assuming a source distance of 10 kpc, together with the sensitivity curves of the advanced LIGO (aLIGO), advanced VIRGO

(AdV), and KAGRA (Abbott et al. 2018). The common feature seen among the models is a spectral peak appearing in a relatively low frequency range of $\sim 100 - 200$ Hz. From the spectrogram in Fig. 9, such low frequency GWs are emitted mainly by convection motions in the postbounce phase. Our most massive progenitor model $z70$ shows a relatively lower peak frequency compared to the other models appearing at $50 \lesssim f \lesssim 100$ Hz. This can be explained by the less compact PNS for model $z70$, although the PNS soon becomes more compact than the rest of models due to a higher mass accretion rate.

Once the initial convection motion ceases, the GW frequency increases with time. While $z9.6$ has only one spectral peak at $f \sim 150$ Hz, due to a considerable subsidence of GWs after $t_{\text{pb}} \gtrsim 100$ ms, the other models exhibit a broadband emission ($s11.2$), extending from ~ 200 Hz to $\gtrsim 1000$ Hz, or rather a double-peaked feature ($s50$ and $z70$). The high frequency emission seen in model $s11.2$ has nearly a comparable amplitude $h_{\text{char}} \sim 2 \times 10^{-21}$ with that of the low frequency emission and is mainly originated from the contracting PNS and its convection. $s50$ and $z70$ also have a high frequency spectral peak appearing at ~ 1700 Hz. From Fig. 9, such a high frequency emission occurs during the second collapse and is associated with the hadron-quark phase transition and its subsequent convection motions as already explained (see also Fig. 8). The amplitude at $f \gtrsim 1000$ Hz in model $s50$ is well above the sensitivity curves of the current-generation GW detectors and could be detectable, if the source locates within our Galaxy.

The present GW analysis for $s50$ is in qualitative agreement with the results reported by Zha et al. (2020), in particular the location of the secondary peaks in excess of 1000 Hz which are related with the dynamics of hadron-quark phase transition. However, the associated magnitude reported here exceeds the one of Zha et al. (2020) by one order of magnitude. It is attribute to the different hybrid EOS, in particular

the significantly lower onset density for the phase transition, below nuclear saturation density, used in the simulation of Zha et al. (2020). This results in a gradual softening of the EOS already at core bounce and hence leading to the continuous contraction of the PNS, lasting for several hundreds of milliseconds post bounce, before the PNS collapses. In our case the EOS softens suddenly, when the PNS enters the hadron-quark mixed phase, which results in the abrupt PNS collapse and a strong GW emission with high amplitudes.

Fig. 12 presents the neutrino luminosities and mean energies for all models². All quantities are evaluated at a distance of $r = 400$ km. Comparing with full GR 2D models in Müller & Janka (2014), the less massive progenitor models $z9.6$ and $s11.2$ show very similar neutrino profiles. We note that the detailed neutrino opacities are different between this study and theirs. The neutrino luminosities of all flavors overlap each other at $t_{\text{pb}} \gtrsim 200$ ms in model $z9.6$. The neutrino mean energies show only a slight increase with time, which indicates that the PNS core does not shrink significantly in this model. At $t_{\text{pb}} = 400$ ms, $L_{\nu_e, \bar{\nu}_e, \nu_x}$ become $\sim 1.2 \times 10^{52}$ erg s⁻¹ and mean energies reach $\varepsilon_{\nu} \sim 11, 14,$ and 16 MeV for $\nu_e, \bar{\nu}_e,$ and ν_x , respectively. In $s11.2$, the accretion luminosities peak at around $t_{\text{pb}} \sim 90$ and reach $L_{\nu_e, \bar{\nu}_e} \sim 4 \times 10^{52}$ erg s⁻¹ and $L_{\nu_x} \sim 2.5 \times 10^{52}$ erg s⁻¹. Afterward they show a decreasing trend. The mean neutrino energies become higher with time and $\varepsilon_{\bar{\nu}_e}$ reaches nearly the same value with ε_{ν_x} at $t_{\text{pb}} \gtrsim 300$ ms. All of these features seen in both $z9.6$ and $s11.2$ are roughly in good agreement with those in Müller & Janka (2014).

$s50$ shows a quite remarkable feature in association with the hadron-quark phase transition. After the first bounce, L_{ν_e} and $L_{\bar{\nu}_e}$ become almost the same values at $t_{\text{pb}} \gtrsim 60$ ms, while L_{ν_x} shows a value of nearly half of the electron type (anti-)neutrino luminosities. When the second collapse initiates, the neutrino luminosities of all flavors drastically increase. From inset panel, we see that $L_{\bar{\nu}_e}$ and L_{ν_e} are showing the highest and lowest luminosity, respectively, and such a hierarchy is consistent with Fischer et al. (2020a). The hierarchy can be explained by the neutron rich environment, through which the shock propagates. As we have explained in Fig. 5, the shock is generated beneath the Y_e -trough and propagates through the neutron rich environment. Therefore abundant electron type antineutrinos are produced via the inverse beta decay process, leading to the highest $L_{\bar{\nu}_e}$. The peak luminosity of $L_{\bar{\nu}_e} \sim 6 \times 10^{53}$ erg s⁻¹ at the second bounce is nearly the same order of that of $L_{\nu_e} \sim 5 \times 10^{53}$ erg s⁻¹ at the first bounce. Such a feature is also reported in Fischer et al. (2020a).

The mean neutrino energies in our most compact model $z70$ show the fastest increase. This is due to the most rapid PNS contraction among the models and to a hottest PNS core as a consequence (Liebendörfer et al. 2004; Sumiyoshi et al. 2007). Shibagaki et al. (2021) reported the neutrino luminosities for the same $70 M_{\odot}$ progenitor model with different nuclear EOS and neutrino opacities. Compared to their results, the neutrino luminosities show slightly higher values of $\sim 20\%$ in all flavors. The neutrino signal associated with the QCD phase transition is not plotted in the figure. This is because we evaluate the neutrino flux at $r = 400$ km and the simulation stops before the neutrino burst reaches that sphere.

Finally we discuss the detectability of the second neutrino burst in association with the hadron-quark phase transition for model $s50$. Fig. 13 illustrates the neutrino detection rate of IceCube (top panel) and HK (bottom panel) as a function of post bounce time for model $s50$ observed at 10 kpc. We evaluate the event rates of each model for two neutrino detectors: IceCube (Abbasi et al. 2011; Salathe et al. 2012) and Hyper-Kamiokande (Abe et al. (2011); Hyper-Kamiokande Proto-Collaboration et al. (2018), hereafter HK). The main channel in these detectors is of anti-electron neutrino ($\bar{\nu}_e$) with inverse-beta decay. For simplicity, we take into account only this channel to evaluate the event rates. We assume that the neutrino energy spectrum is a Fermi-Dirac distribution (see Lund et al. (2010); Takiwaki & Kotake (2018) for more detail). Furthermore in the evaluation, we consider two extreme cases: all $\bar{\nu}_e$ emitted from the source reach the detectors without neutrino flavor conversion and cause the signal at the detectors (red lines in the figure); all $\bar{\nu}_x$ (identical to ν_x in this study) emitted from the source are completely swapped by $\bar{\nu}_e$ and cause the signals (black lines).

From the plot, we find that, during the first ~ 50 ms after bounce, the $\bar{\nu}_x$ signal shows a larger detection rate for both detectors. Afterward the $\bar{\nu}_e$ signal becomes stronger than that of $\bar{\nu}_x$. After $t_{\text{pb}} \sim 100$ ms, both signals show nearly constant values until the onset of the second collapse. Such features are consistent with previous normal CCSN models (Tamborra et al. 2013; Nagakura et al. 2021b). Once the second collapse takes place, the signals become approximately 30 times stronger than the values before the second collapse. The peak values for both $\bar{\nu}_e$ and $\bar{\nu}_x$ signals reach nearly the same value of ~ 5000 ms⁻¹ (IceCube) and ~ 700 ms⁻¹ (HK). Fischer et al. (2018) reported the number of expected neutrino events in the Super-Kamiokande (SK) detector of ~ 100 with 2 ms bins from the second neutrino burst event, assuming a source distance of 10 kpc. Concerning the fiducial volume of HK, which is approximately 10 times larger than that of SK, our expected event rate ~ 700 ms⁻¹ in HK can be converted to ~ 140 signals with 2 ms bins, which is in line with that of Fischer et al. (2018).

² We have checked the reliability of current neutrino transport method in the Cartoon coordinates using the standard neutrino opacity set, which is summarized in App. B.

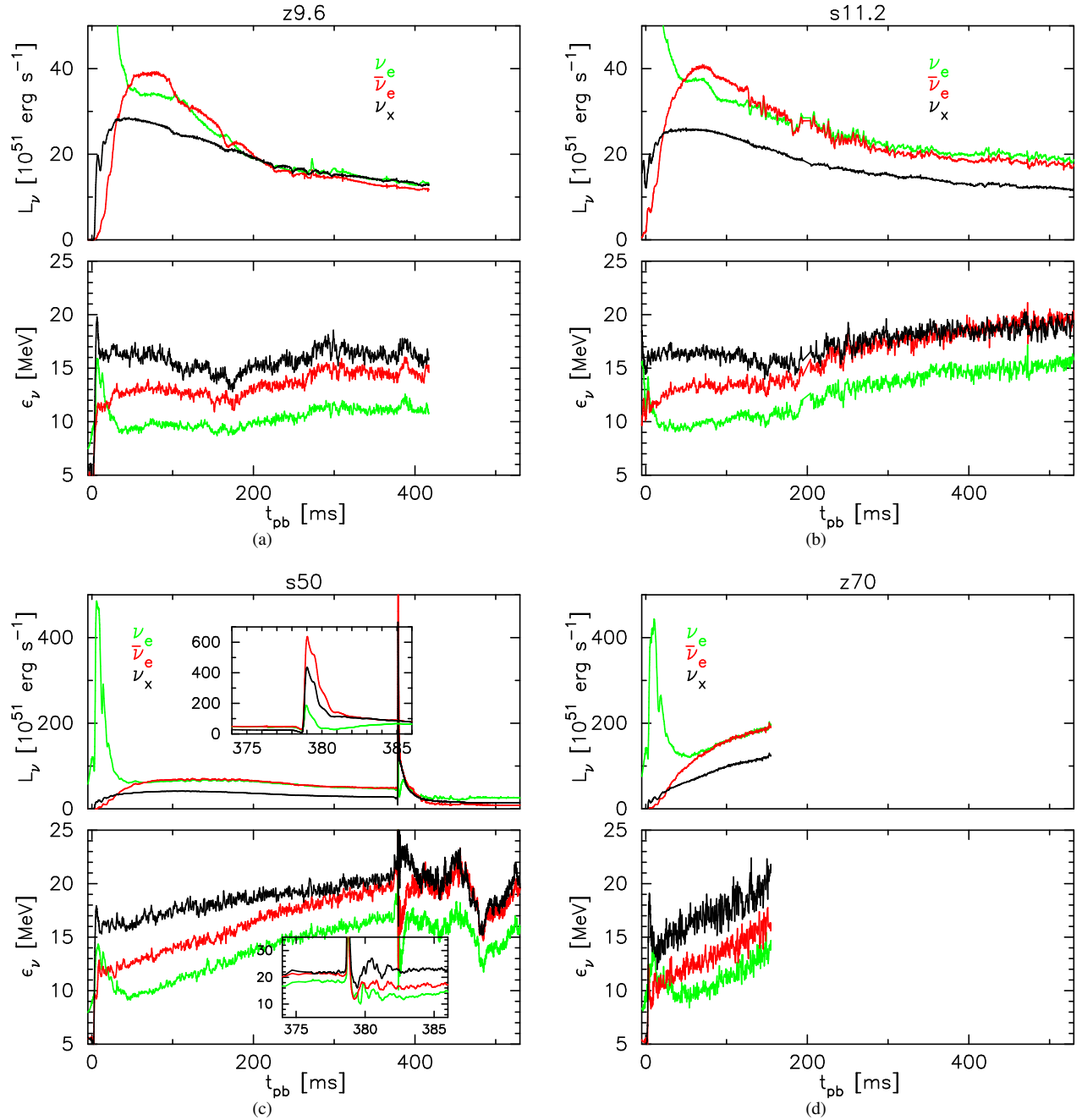


Figure 12. Post-bounce evolution of the neutrino luminosities (top panels) and mean neutrino energies (bottom panels) for all models. Green, red, and black lines belong to ν_e , $\bar{\nu}_e$, and ν_x , respectively. For *s50*, the inlays show the second burst.

5. SUMMARY AND DISCUSSION

In this study, we have presented possible effects of the phase transition from hadronic matter to quark matter in CCSNe. For this purpose, we have performed 2D axisymmetric numerical simulations of core collapse of massive stars with $9.6 M_\odot$, $11.2 M_\odot$, $50 M_\odot$, and $70 M_\odot$. Our core-collapse supernova model is based on numerical relativity, which solves the general relativistic neutrino radiation hydrodynamics equations together with the two-moment (M1) neutrino transport

equations of Kuroda et al. (2016b). We use up-to-date neutrino opacities following Kotake et al. (2018). We employed the hybrid EOS DD2F-RDF 1.2 of Fischer et al. (2018) and Bastian (2021), which considers a first order hadron-quark phase transition.

Our results demonstrate a diversity of post bounce evolution depending on the progenitor. We confirm that low and intermediate mass progenitors, examined at the examples of a low-metallicity model with ZAMS mass of $9.6 M_\odot$ and a solar-metallicity model with ZAMS mass of $11.2 M_\odot$, belong

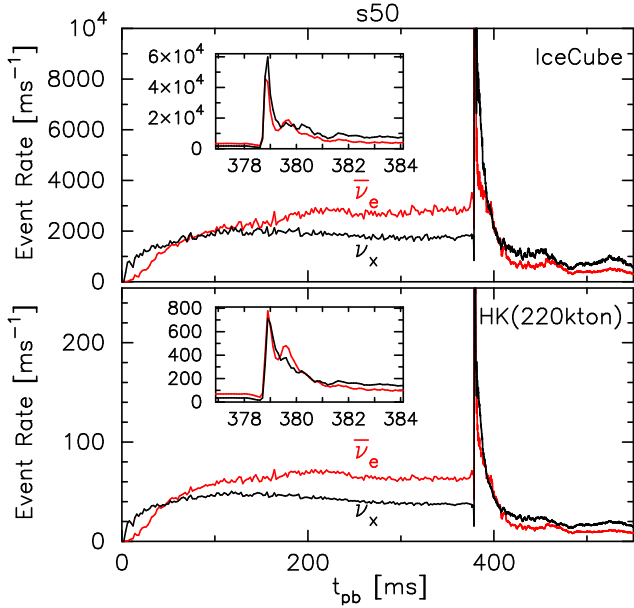


Figure 13. Neutrino detection rate of IceCube (top panel) and HK (bottom panel) as a function of time after bounce for model $s50$ observed at 10 kpc. The red and black lines are for $\bar{\nu}_e$, assuming no flavor conversion, and for $\bar{\nu}_x$, assuming complete flavor conversion, respectively. The inlays show the second burst.

to the class of core-collapse supernova explosion driven by the neutrino-heating mechanism. The latter one is aided by the presence and development of convection. Both neutrino-driven explosion models are in qualitative agreement with previous studies, e.g., in terms of the explosion dynamics as well as the diagnostic explosion energy obtained. None of these models reach the conditions required for the hadron-quark phase transition, prior to the onset of the supernova explosion. In contrast, the more massive progenitors explored with ZAMS mass of $50 M_{\odot}$ of solar metallicity and an extremely metal-poor star with ZAMS mass of $70 M_{\odot}$, both undergo the phase transition during the post-bounce mass accretion phase. However, an explosion is only obtained for the $50 M_{\odot}$ model while for the $70 M_{\odot}$ model we observe a failed explosion and black-hole formation instead. This is related to the maximum mass of the DD2F-RDF 1.2 hybrid EOS, which the latter one exceeds at the onset of the hadron-quark phase transition. Contrary therefore, for the $50 M_{\odot}$ model we report the formation and propagation of a second shock wave, as a direct consequence of the hadron-quark phase transition, which determines the onset of the supernova explosion, in qualitative agreement with the results reported by Fischer et al. (2018) and Fischer et al. (2020b). Nevertheless, the yet incompletely understood question about the dependence on metallicity, preliminarily addressed in Fischer (2021), remains to be explored in a more systematic study.

Our finding inherent to the multi-D nature is the strong RT mixing. In the context of CCSNe, the potential re-

gion of development of the RT mixing is usually distant from the PNS, e.g. the interface between CO core and He layer (Wongwathanarat et al. 2015), which locates at $r \gtrsim 10^{9-10}$ cm depending on the progenitor model. However, because of the shallow density profile in $s50$, the RT mixing can develop even at a distance of $r \sim 10^8$ cm. Because of the short distance, the RT mixing can develop in the early post explosion phase of a few 10 ms after the explosion onset has been launched in terms of the initial propagation of the second shock wave. A combination of the PDE and a prompt development of the RT mixing might have an influence on the nucleosynthetic yields previously reported in Fischer et al. (2020b).

Furthermore, we discussed possible effects of the hadron-quark phase transition on the multi-messenger observables, gravitational waves and neutrinos. For the less massive progenitor models $s9.6$ and $s11.2$, we confirm commonly found features in their GW signals: $z9.6$ shows a subsidence of GW emission, once it entered a runaway explosion phase at $t_{\text{pb}} \gtrsim 200$ ms, and $s11.2$ exhibited a rather constant GW emission even after its explosion, though the peak frequency increases with time in association with the PNS contraction. The massive progenitor stars $s50$ and $z70$ display a quiescent phase for GW emission once their standing shock front starts to recede. However, in full 3D models, the SASI phase could arise in this period and contribute to the GW emission. The contribution can be significant to produce sizeable GWs according to the previous 3D studies (e.g., Kuroda et al. 2016a). At the second bounce in $s50$, ring down of the quark core produces strong and high-frequency GWs, with their amplitudes and typical frequencies exceeding $Dh \sim 250$ cm and 1 kHz, respectively. Since the Brunt-Väisälä frequencies behind the second bounce shock are reaching a few kHz, the high-frequency GWs could be an outcome of these convection motions (see also the supplemental material of Zha et al. 2020, for the relation between the convection motions behind the second bounce shock and typical GW frequencies.). Although their typical frequency appears at kHz range, which is out of the best sensitivity of current ground-based gravitational wave detectors, their signal is strong enough to be detected, if the source locates at 10 kpc. Even in the failed explosion case $z70$, the region outside the apparent horizon contributes significant GW emission. Thus, we can expect that the signal of a hadron-quark phase transition might be imprinted in GWs even in the failed CCSNe, although the detailed waveform should be examined in full 3D models, considering the SASI development.

Concerning the neutrino profiles, the neutrino luminosity and mean energy of all flavors have shown quite reasonable values in models $z9.6$ and $s11.2$ with previous studies (e.g. with 2D relativistic models of Müller & Janka 2014), since these two less massive progenitor models do not undergo the

hadron-quark phase transition during our simulation time. While in $s50$, we have observed the second neutrino burst when the second core-bounce takes place. In contrast to the first bounce, the electron type antineutrino luminosity $L_{\bar{\nu}_e}$ exhibits the highest luminosity among the flavors. This is because the quark core bounce occurs below the Y_e -trough, and the shock propagates through the neutron rich environment. Therefore, the inverse beta decay becomes the dominant reaction process and produces numerous electron type antineutrinos. The hierarchy seen in the second neutrino burst is consistent with previous studies (Fischer et al. 2018; Zha et al. 2020).

Recently long-time 2D SN simulations, longer than several seconds in 2D (c.f. Müller 2015; Nakamura et al. 2019; Burrows & Vartanyan 2021; Nagakura et al. 2021a) and even in 3D by Bollig et al. (2021) become computationally feasible. Because of our short simulation times, the diagnostic explosion energy of models $z9.6$ and $s11.2$ do not reach energy saturation. According to Müller (2015), the energy saturation can be reached after ~ 5 s for the model $s11.2$. Such long-time simulations are necessary to discuss the diagnostic explosion energy, as well as the final property of remnant star. It is worth mentioning that the inclusion of progenitor’s precollapse inhomogeneities in the burning shells (Yadav et al. 2020; Yoshida et al. 2019, 2020, 2021) has been pointed out to make the underpowered explosion more en-

ergetic (Müller & Janka 2014; Bollig et al. 2021). Starting from the multi-D progenitors, we have to study how the multi-D hydrodynamic motions in the pre-explosion phase result in the multi-messenger signals up to the late explosion phase (see previous attempts in Nakamura et al. (2019), Nagakura et al. (2021a), and Witt et al. (2021)). Though this is not an easy task, this should be the only promising way we should be on to enhance the predictive power of CCSN multi-messenger signals as quantitatively as possible.

ACKNOWLEDGMENTS

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APPENDIX

A. SPATIAL RESOLUTION TEST

In this appendix, we are showing spatial resolution dependence of the post-bounce PNS evolution to explain that the currently used spatial resolution is adequate to follow the quark core evolution. For that purpose, we perform three CCSN simulations “case1-3” with different spatial resolution. “case1” is identical to the model $s50$ discussed in the main text and employs a nested structure with the maximum refinement level of $L_{\max} = 10$, if the maximum rest mass density ρ_{\max} is less than $8 \times 10^{14} \text{ g cm}^{-3}$, and with $L_{\max} = 11$, once the maximum rest mass density exceeds $8 \times 10^{14} \text{ g cm}^{-3}$. Consequently the central numerical grid width achieves ~ 115 m when $\rho_{\max} > 8 \times 10^{14} \text{ g cm}^{-3}$ (or ~ 230 m when $\rho_{\max} < 8 \times 10^{14} \text{ g cm}^{-3}$). While in “case2”, we set the maximum refinement level to be $L_{\max} = 10$, i.e. the finest numerical grid width is ~ 230 m throughout the calculation. This models is to check the resolution dependence after the quark core formation and also to see whether the quark core in current model $s50$ is adequately resolved. “case3” is calculated to see the resolution dependence particularly focusing on the PNS contraction phase and use $L_{\max} = 9$, i.e. the finest numerical grid width is ~ 450 m.

In the left panel in Fig. 14, the maximum density evolution is plot for three cases “case1-3”. In the right panel, we show the maximum density (top) and ADM mass (bottom) for case1 (red line) and case2 (light blue), focusing on the post second bounce phase. We first compare the case1 (red) and case3 (black) to check the impact of numerical resolution on the second collapse time. The finest numerical grid in these two models is ~ 230 m (red) or ~ 450 m (black) during the PNS contraction phase ($t_{\text{pb}} \lesssim 350$ ms). From the left panel, the black line systematically shows a smaller maximum density compared to that of case1. Although the difference already appears at just after the first bounce, we also find that the central entropy secularly increases in case3 (not shown in the figure). This is due to the numerical dissipation in lower resolution model case3. Therefore, the maximum density in case3 is showing smaller value throughout the PNS contraction phase. However the second collapse times in case1 and case3 are exhibiting a rough convergence, only ~ 20 ms difference between red and black lines. We can thus say that, for the PNS contraction phase, a rough numerical convergence is achieved for grid width smaller than ~ 450 m.

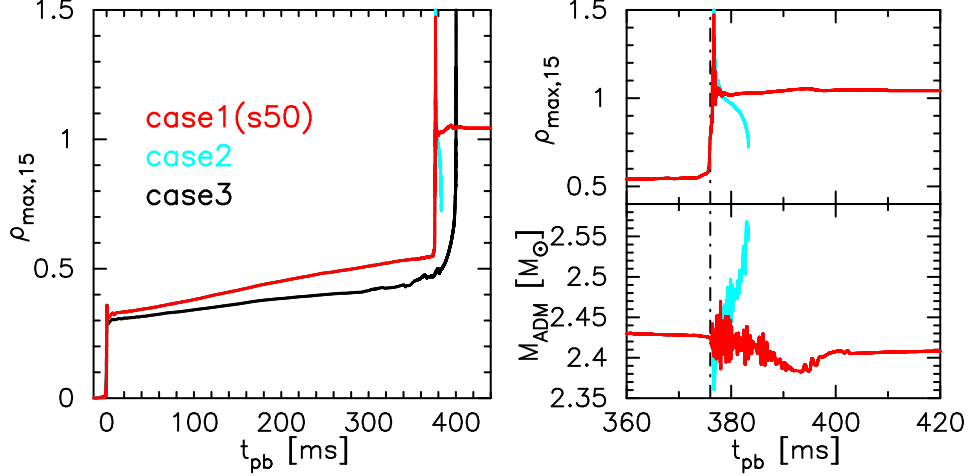


Figure 14. *Left panel:* Evolution of the maximum density $\rho_{\max,15}$ for three cases “case1-3” (see text for their numerical setup). *Right panel:* Zoom in for the evolution of the maximum density $\rho_{\max,15}$ (top panel) and ADM mass M_{ADM} (bottom panel) for case1 (red line) and case2 (light blue), focusing on the post second bounce phase. The vertical dot-dashed lines indicate when we raise the maximum refinement level from $L_{\max} = 10$ to 11.

Next we compare case1 and case2. From top right panel, it is obvious that the maximum density drastically decreases after the second core bounce in case2 (light blue). At the same time, the ADM mass significantly increases and does not settle into a new quark core branch as case1 does. This is because the lower resolution model case2 suffers from a severe numerical dissipation and the quark core seems to be artificially gravitationally unbound. From these comparisons, we conclude that a resolution of ~ 110 m is required to adequately resolve the quark core evolution.

B. TEST FOR THE NEUTRINO TRANSPORT IN A MIXED FRAME OF CYLINDRICAL AND CARTOON COORDINATES

This appendix is devoted to test our neutrino transport in a mixed frame of cylindrical and Cartoon coordinates, in which all hydrodynamical and neutrino-radiation quantities, as well as the metric variables, are not evolved straightforwardly as in the Cartesian coordinates used in our previous studies (e.g., Kuroda et al. 2016b). To this end, we perform the same CCSN simulation reported by O’Connor et al. (2018) and follow the collapse of the solar metallicity, non-rotating $20 M_{\odot}$ star taken from Woosley & Heger (2007), and compare the neutrino profiles. Furthermore, the neutrino opacity is also the same with O’Connor et al. (2018) and uses the baseline set of Bruenn (1985) including Brehmsstrahlung of Hannestad & Raffelt (1998). The energy bins are 20 bins, which logarithmically cover the region of 0-300 MeV. One of the major difference from O’Connor et al. (2018) is the spherical symmetric property, which we cannot take into account in the Cartoon coordinates based numerical code. We also note that the presented model is calculated by numerical relativity and thus might have some inherent differences from those in O’Connor et al. (2018). Therefore, in the late post bounce phase, when the convection motions become dominant, our neutrino profiles show different results from those in O’Connor et al. (2018). In Fig. 15, we plot neutrino luminosities (left) and mean neutrino energies (right) extracted at a radius of $r = 400$ km.

We first explain comparison of neutrino luminosities. From left panel, the accretion luminosity reaches $\sim 7.5 \times 10^{52}$ erg s $^{-1}$ for ν_e and $\bar{\nu}_e$ and $\sim 4.1 \times 10^{52}$ erg s $^{-1}$ for ν_x . The post bounce time, when those accretion luminosity reach their maximum, as well as their peak values are in good agreement with O’Connor et al. (2018). At around $t_{\text{pb}} \sim 240$ ms, the Si-O interface accretes and the mass accretion rate suddenly drops. Consequently $L_{\nu_e, \bar{\nu}_e}$ significantly decreases. Afterward it enters the convection dominant phase and the neutrino profile in multi-D model starts to diverge from that in spherically symmetric model. Indeed due to the vigorous convection motions, our neutrino luminosities show systematically higher values in all flavors than those in 1D models of O’Connor et al. (2018), roughly 25% and 100% increase for $L_{\nu_e, \bar{\nu}_e}$ and L_{ν_x} , respectively, at $t_{\text{pb}} = 400$ ms. Such a higher neutrino luminosity in multi-D model reflects the effects of PNS convection around the neutrino sphere (Buras et al. 2006).

Regarding the neutrino mean energy (right panel), we again find a consistent result with O’Connor et al. (2018), at least for the first ~ 100 ms after bounce. During that early epoch, $\varepsilon_{\bar{\nu}_e}$ reaches its minimum value of ~ 9 MeV at $t_{\text{pb}} \sim 50$ ms. $\varepsilon_{\bar{\nu}_e}$ is showing a higher value than ε_{ν_e} , preserving the deviation $\varepsilon_{\bar{\nu}_e} - \varepsilon_{\nu_e}$ to be ~ 3 MeV throughout most of the post bounce phase. This feature is also seen in O’Connor et al. (2018). While ε_{ν_x} stays at ~ 17 MeV just after bounce $t_{\text{pb}} \lesssim 100$ ms, and then gradually increases, though the increase rate is not so high as those of ε_{ν_e} and $\varepsilon_{\bar{\nu}_e}$. Due to the multi-D effects, all the mean neutrino energies become

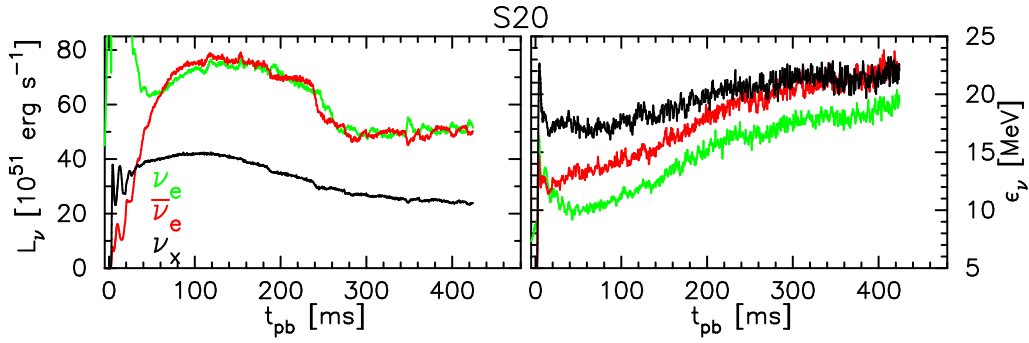


Figure 15. Post-bounce evolution of the neutrino luminosities (left panel) and mean neutrino energies (right panel) for the solar metallicity, non-rotating $20 M_{\odot}$ progenitor star from Woosley & Heger (2007). The latter is identical to the one used in a global comparison of CCSNe simulations reported by O’Connor et al. (2018). The neutrino opacity is the same as was used in O’Connor et al. (2018). Green, red, and black lines belong to ν_e , $\bar{\nu}_e$, and ν_x , respectively.

systematically higher than the 1D results, especially in the late post bounce phase $t_{\text{pb}} \gtrsim 200$ ms. From these comparisons, we conclude that the new neutrino transport method in a mixed frame of the cylindrical and Cartoon coordinates works consistently with other numerical codes and is capable to perform CCSN simulations.

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